II. NUCLEAR STRUCTURE AND DYNAMICS

CONTENTS

INTRODUCTION

Investigations of nuclear theory community at BLTP cover a large part of contemporary nuclear physics. Nuclear theorists published 157 papers in peer reviewed journals in 2009-2010. Moreover, they gave talks at several dozens international conferences, schools and workshops over the world.
The whole area of nuclear physics studies at BLTP was divided (somewhat formally) into four projects

- Nuclear Structure Far from the Valley of Stability
- Nucleus-Nucleus Interactions and Nuclear Properties at Low Excitation Energies
- Exotic Few-Body Systems
- Nuclear Structure and Dynamics at Relativistic Energies

The first two projects dealt with the low-energy nuclear physics. The activity in this direction has the longest tradition at BLTP. In the present report it is represented by 13 contributions.
In the first contribution by S. Mishev and V. V. Voronov the extended version of the Quasiparticle-Phonon Model for odd spherical nuclei is presented. This is the essentially nonlinear approach consistently treating new types of ground state correlations in both the even-even core and odd nucleus itself. In two subsequent notes (by R. V. Jolos and P. von Brentano (University of Cologne) and N. Yu. Shirikova *et al.*) particular ways to improve the description of so-called "soft" nuclei are presented.
In the following three contributions some aspects of the theory of nuclear giant resonances are considered. A. P. Severyukhin *et al.* shortly summarize the results obtained in 2009-2010 employing a finite rank separable approximation for the residual interaction within the Skyrme-Quasiparticle-Random-Phase approach (QRPA). This approximation developed several years ago in principle allows one to go beyond the self-consistent Skyrme HFB-QRPA and take into account the coupling of one-phonon QRPA states with numerous two-phonon states. In the contribution, calculations with the phonon-phonon coupling are presented for the isoscalar $E2$ resonance in ^{132}Sn. In the contribution by V. O. Nesterenko *et al.* it is shown that the attempt to describe available data on the spin-flip $M1$ resonance within the standard Skyrme HFB-QRPA approach meets some problems, thus demanding a thorough revision of spin-dependent terms in the Skyrme functional. R. G. Nazmitdinov discusses a general approach to characterize fluctuations of measured cross sections of nuclear giant resonances arguing that the spreading width of a nuclear resonance is determined by the number of fragmentations over more complicated configurations.
In the contribution by F. Šimkovic *et al.* the new achievements in calculations of nuclear matrix elements of neutrinoless double beta ($0\nu\beta\beta$) decay are presented. This international group has strongly contributed to the subject. Although predictions of different nuclear models still noticeably diverge, the presented results are encouraging. The second contribution dealing with weak processes in atomic nuclei is the one by A. A. Dzhioev *et al.* Here the main accent is put on weak processes in nuclei at finite temperatures with applications to astrophysical problems.

Among other problems of nuclear structure physics which were investigated by BLTP nuclear theorists during these two years but are left beyond the scope of the report we would like to mention the studies of low-lying excited states in deformed nuclei, especially isomers in very heavy nuclei, and of multiple reflection-asymmetric type band structures in nuclides of the actinide region.

The next five contributions are devoted to problems of nuclear reaction theory. In the three of them the dinuclear system (DNS) model is exploited. In the first contribution by A. S. Zubov *et al.* a model of formation of hyperdeformed states in the entrance channel of heavy-ion reaction is proposed. The authors determine the optimal reactions and conditions (bombarding energies, range of angular momenta) for the identification of hyperdeformed states. It should be noted that these studies are supported by the grant of the President of the Russian Federation for young PhD scientists. G. G. Adamian *et al.* apply the DNS model to analyze the production of very neutron-rich nuclides ^{84,86}Zn and ^{90,92}Ge in the multinucleon transfer actinide-based reactions with a ^{48}Ca beam with existing beams and detection systems. A. K. Nasirov *et al.* exploit the DNS model to distinguish contributions of the quasifission and fast fission processes to the cross sections of selected nucleus-nucleus collisions.

The next two contributions are connected with the experimental program of the Flerov Laboratory of Nuclear Reactions at JINR. S. N. Ershov presents selected results of recently developed microscopic four-body distorted wave theory for two-neutron halo breakup reactions leading to low-lying halo excitations which accounts for both elastic and inelastic breakup. V. K. Lukyanov *et al.* analyzes the ^{6}He+^{12}C elastic scattering data with the developed microscopic optical potential and discusses its advantages and shortcomings.

Subjects of investigations within the few-body theory field are rather scattered. They include a consistent three-body treating of the Helium trimer systems, further developments and new applications of the Dubna-Mainz-Taipei meson-exchange model, dynamics of resonant molecule formation in waveguides, cluster description of the famous Hoyle state in ^{12}C, sharp norm bounds on variation of invariant subspaces for multi-channel Hamiltonians, rigorous mathematical results in the theory of three-body collisions. In the report one can find the contributions covering the three of the listed topics.

V. V. Pupyshev discusses the new peculiarities of the three-body wave functions near the triple impact point found by him recently. The investigations by V. S. Melezhik and his collaborators are devoted to few-body physics at low dimensions. They found and studied two novel effects in the ultracold atomic collisions in harmonic traps. The important point is that the theoretical predictions impacted experimental efforts and were confirmed by them. The interesting subject to study is the interaction of nucleons and mesons consisting of different quarks (e.g., strange quarks) since they can exchange mainly by gluons. Just such systems $\phi + 2N$ and $\phi + 3N$ are considered in the contribution by V. B. Belyaev *et al.*.

The last five contributions cover the area of nuclear physics at relativistic energies. V. K. Lukyanov *et al.* constructs in the high-energy approximation the microscopic optical potential for the K^{+}-nucleus elastic scattering. New achievements of the Bethe-Salpeter approach is presented in the contribution by S. G. Bondarenko, V. V. Burov and E. P. Rogochaya. The authors construct the new rank-6 separable interaction kernel and apply it to describe various characteristics of the deuteron. A thorough comparison with other approaches is performed as well. V. D. Toneev and A. S. Khvorostukhin investigate the bulk and shear viscosity of hadron matter within the elaborated earlier relativistic mean-field model with scaled hadron masses and couplings extended to finite temperatures. L. P. Kaptari and C. Ciofi degli Atti (INFN, Perugia) analyze the reaction $^3\mathrm{He}(e, e'p)pn$ within a parameter-free approach based upon realistic few-body wave functions and treating the rescattering of the struck nucleon. The subject of the contribution by A. I. Titov is somewhat unusual for the community involved in pure theoretical studies. Here it is estimated whether the high-energy laser-driven electrons can produce a sizeable amount of muon pairs. The positive answer means that the produced $\mu^{\mp}$ can be used in studying various aspects of muon-related physics in tabletop installations.

Certainly, being quite wide the activity area of BLTP nuclear theorists does not cover the whole field of nuclear physics. The reader can notice, however, that the presented topics mainly reflect current trends in this branch of exact sciences.

V. V. Voronov

A. I. Vdovin

GROUND-STATE CORRELATIONS AND STRUCTURE OF THE LOW-LYING STATES IN ODD-EVEN SPHERICAL AND TRANSITIONAL NUCLEI

S. Mishev and V. V. Voronov

The influence of the Pauli principle and the nucleon-nucleon correlations in the ground states of spherical and transitional even-even nuclei on the structure of the low-lying states in odd-even nuclei was examined in [1, 2]. We studied correlations caused by the quasiparticle-phonon interaction in the ground state beyond the pairing correlations. The effects owing to the ground-state correlations (GSC) are becoming essential as the number of nucleons in the unclosed shells increases. Two aspects of this problem were addressed. The first considered was the suggestion that quasiparticle and quasiparticle⊗phonon states could exist in the ground states of even-even nuclei. By analogy with the random phase approximation (RPA) for even-even nuclei this implied non-vanishing backward amplitudes in the odd-A nucleus wave function. In earlier studies related to this subject, the quasiparticle and phonon operators were taken as commuting ones, thus neglecting the Pauli principle which can be unsatisfactory in a number of nuclei, since in them serious deviations from the independent harmonic motion occur. In these cases the disregard of the innate fermion structure of the phonons is unjustified. In Ref. [1], we performed analytical calculations following the exact commutation relations between the quasiparticle and phonon operators and evaluated the effects of the resulting corrections on the spectra and single-particle spectroscopic strengths of states in the vicinity of the Fermi level in a number of odd-even barium isotopes. It was found that the first and second states with the same angular momentum and parity become closer in energy than the predictions of models disregarding the backward amplitudes which turned out to be in accord with the experimental data. A considerable shift of the single-particle fragmentation to higher energies influenced by the GSC was also registered. In this treatment the Pauli exclusion principle manifested itself in the emergence of factors $(1-\mathcal{L}(Jj\lambda i))$ which turn to zero whenever a particular three-quasiparticle state is disallowed.

The second aspect of this problem is associated with an improved method for calculating average values of various quantum-mechanical operators, which is consistent with the concept that the nuclear ground state is correlated. This method, which is a generalization of the quasiparticle RPA (QRPA) and referred to as the Extended RPA (ERPA), was suggested by Hara [3] and Ikeda *et. al.* [4]. The approach, they have proposed, broadens the area of applicability of the conventional theory which relies on the assumption that the true ground state must not be very different from the quasiparticle vacuum state. In ERPA the quasiparticle occupation numbers enter into the basic equations of the theory explicitly, leading to a codependence between the different layers of the theory otherwise separated in QRPA. Further developed [5] and applied to concrete nuclei [6], this approach proved successful in improving the theoretical results for most measurable quantities near the nuclear ground states as, for example, the transition charge densities in the interior region. Following the ERPA prescription, we derived [2] renormalized expressions for the interaction between quasiparticles and phonons in both the ground and excited states. This interaction depends on the quasiparticle occupation numbers in the ground state explicitly, thus coupling the core-particle equations with the generalized equations describing the pairing correlations and the excited vibrational states of the even-even core forming a large nonlinear system. The superiority of ERPA over QRPA in reproducing the experimental data on $B(E2|\text{g.s.} \to 2_1^+)$ in even-even transitional nuclei stimulated a survey on the effects of the GSC on the electric transition probabilities in odd-even systems. Despite the considerable enhancement in these quantities due to the GSC, it was concluded that further correlation effects need to be taken into account in order to reach better agreement with the experimentally measured values. Numerical calculations on the spectroscopic factors in several Te, Xe, and Ba isotopes were also performed, indicating an overall improved description due to the weakened quasiparticle-phonon interaction strengths in the renormalized version of the model.

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X(5)* MODEL AND $N = 90$ NUCLEI

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In the collective nuclear models using the Bohr Hamiltonian the mass tensor plays as an important role as the potential energy. However, the effect of the mass tensor on the collective motion is not seen as easily as the effect of the potential energy. The shape of the potential tells us whether the nucleus is spherical, deformed or transitional. This explains partly why in many publications the kinetic energy term of the collective quadrupole Hamiltonian is taken in its simplest form. This simplest form assumes a constant mass coefficient if Bohr's usual shape variables are used. However, our analysis of the experimental data has shown [1] that if, as usual, the nuclear shape variables $\alpha_{2\mu}$ which are proportional to the quadrupole moment operator are used as the collective variables in the Hamiltonian then the mass tensor should have a complicated form. It will have not only monopole but also quadrupole and hexadecupole components and therefore cannot be reduced to one constant mass coefficient. This conclusion was made basing on the analysis of the sum rules calculated as the matrix elements of a double commutator of the quadrupole moment operator with a Hamiltonian [1].

Table 1: The calculated values of the B(E2)'s obtained for the collective quadrupole model with the X(5) Hamiltonian and the modified E2 transition operator. The parameters of the modified E2 transition operator are found by fitting the experimental values of $B(E2; 0^+_\beta \to 2^+_1)$ and $B(E2; 2^+_\gamma \to 2^+_1)$. The strong interband E2 transitions are marked by a dot •. The experimental data for ^{150}Nd, ^{152}Sm and ^{154}Gd are included in the Table for comparison with the calculated results. All values are normalized to the $2^+_1 \to 0^+_1$ transition. β_w is the maximum possible value of β allowed by the X(5) potential which is taken to be equal to 0.70.

$B(E2; I' \to I)$	^{150}Nd exp	^{150}Nd X(5)* $\chi_2\beta_w$= -0.302 $\chi_3\beta_w^2$= -1.410	^{152}Sm exp	^{152}Sm X(5)* $\chi_2\beta_w$= -0.186 $\chi_3\beta_w^2$= -1.671	^{154}Gd exp	^{154}Gd X(5)* $\chi_2\beta_w$= -0.172 $\chi_3\beta_w^2$= -1.477
$2^+_1 \to 0^+_1$	100(2)	100	100(2)	100	100(2)	100
$4^+_1 \to 2^+_1$	158(2)	153	145(2)	149	155(4)	151
$6^+_1 \to 4^+_1$	183(2)	180	170(3)	171	168(4)	175
$8^+_1 \to 6^+_1$	242(22)	197	198(10)	183	197(11)	190
$2^+_\beta \to 0^+_\beta$	99(20)	76	74(19)	75	34(3)	76
• $0^+_\beta \to 2^+_1$	34(2)	34	23(1)	23	28(2)	28
$2^+_\beta \to 0^+_1$	1.0(2)	0.2	0.6(1)	0.003	0.34(3)	0.03
$2^+_\beta \to 2^+_1$	7.8(20)	2.8	3.8(3)	1.1	4.0(3)	1.8
• $2^+_\beta \to 4^+_1$	15(3)	19.5	13(1)	13.1	13(1)	16
• $2^+_\gamma \to 2^+_1$	5.0(2)	5.0	6.5(3)	6.5	7.8(6)	7.8

As an application of these ideas we generalize the so called X(5) model. This model was introduced as a simple model which describes surprisingly well the $N = 90$ transitional

nuclei. Our generalization of the X(5) model consists of keeping the X(5) Hamiltonian but generalizing the E2–operator.
The fact that the mass tensor in the collective Hamiltonian cannot be taken as a constant and should be considered as a function of the collective coordinates, which is the case when the quadrupole and the hexadecupole components are presented, makes it much more difficult to solve the Schrödinger equation. A solution is much simpler if the collective quadrupole variables are chosen in such a way that the mass tensor is reduced to one constant mass coefficient. Therefore, we should consider a more complicated expression for $Q_{2\mu}$ if the mass tensor in the collective Hamiltonian is reduced to a one-constant mass parameter. Following this idea, we suggest for $Q_{2\mu}$ an expression with three parameters, i.e.

$$Q^{II}_{2\mu} = q\,(\tilde{\alpha}_{2\mu} + \chi_2 \cdot (\tilde{\alpha}\tilde{\alpha})_{2\mu} + \chi_3 \cdot \tilde{\alpha}_{2\mu}(\tilde{\alpha}\tilde{\alpha})_0)\,. \tag{1}$$

The results of the calculations [2] of the B(E2) values with this form of the quadrupole moment operator are presented in Table 1.

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$E0$ TRANSITIONAL DENSITY FOR NUCLEI BETWEEN SPHERICAL AND DEFORMED SHAPES

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Electric monopole transitions between the eigenstates of even–even nuclei can provide important information about the properties of low-lying collective nuclear states. The average value of the monopole operator is proportional to the nuclear radius in that particular state. The nondiagonal matrix elements of the $E0$ operator are sensitive to the distribution of the collective wave functions over the axial deformation β. The $E0$ transitional densities contain even more detailed information on the collective wave functions than the mere $E0$ transition strengths. The transitional densities are very sensitive to the dependence of the radial density distributions on deformation. Phenomenological models employ rather smooth functions for description of the nuclear density, and it is interesting to investigate this problem in the framework of microscopical models.

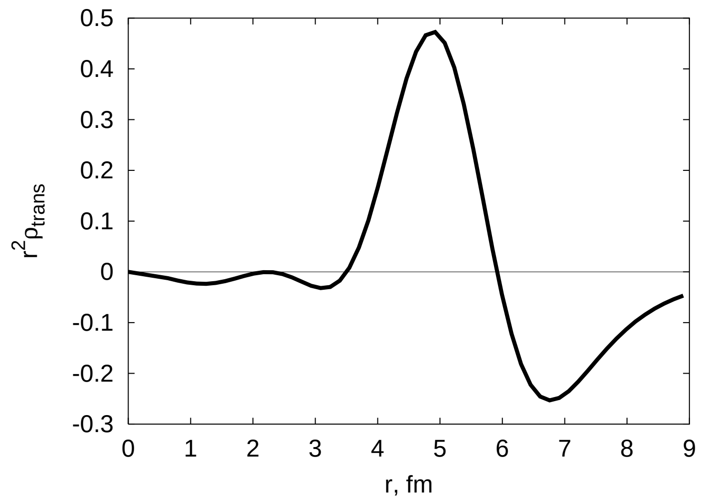

Figure 1: $E0$ transition density for the $0^+_{gs} \to 0^+_2$ transition multiplied by r^2. The calculations are performed for ^{150}Nd. The wave functions of the 0^+_{gs} and 0^+_2 states are the solutions of the Bohr Hamiltonian with the potential corresponding to the X(5) limit.

There are well developed phenomenological, both geometrical and algebraic, models which became very popular in connection with discussions of the shape phase transitions in nuclei. These phenomenological approaches can be used to describe the properties of the transitional nuclei. Their numerical application is quite simple. Of course, the fully microscopic approaches are more appropriate. However, frequently, phenomenological models, which use the experimental data to fix the model parameters, are closer to the experimental results. This situation suggests the possibility to develop the following partially microscopic approach to calculation of the $E0$ transitional densities for nuclei which are characterized by the large amplitudes of the collective quadrupole motion. In order to realize this approach, it is necessary, at first, to construct in the framework of a microscopic model an effective nuclear density operator as a function of the radius, deformation parameter β, and collective momentum $\partial/\partial\beta$. As the second step, the matrix elements of this operator should be calculated between the eigenstates of the phenomenological collective Hamiltonian.

The full expression for the effective $E0$ transitional density operator $\rho_{\text{eff}}(r,\beta)$ was obtained using a technique described in [1]. The result is

$$\rho_{\text{eff}}(r,\beta) = \langle\beta|\rho\hat{(r)}|\beta\rangle + F(r) + \left(\frac{1}{2Z}F(r) - G(r)\right)\frac{\partial^2}{\partial\beta^2}, \tag{1}$$

where $F(r)$ and $G(r)$ are functions whose expressions are given in [1]. Equation (1) has been used to calculate both proton and neutron densities. The results obtained demonstrate that only the first term in (1) is important.

Figure 1 finally displays the results of our calculations of the $0^+_{gs} \to 0^+_2$ $E0$ transitional density for ^{150}Nd, as an example.

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VIBRATIONAL EXCITATIONS AND A SEPARABLE APPROXIMATION FOR SKYRME INTERACTIONS

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Among recent developments in the Skyrme-Quasiparticle-Random-Phase approach (QRPA) a finite rank separable approximation [1,2] for the residual interaction seems to be particularly promising. The method enables one to perform calculations in very large configuration spaces. Moreover, it has been generalized to take into account the coupling between one- and two-phonon configurations in wave functions of excited states [3]. Applications of the method to study the low-lying as well as high-lying modes of various vibrations can be found in Refs. [4-6].

In Ref. [4], the evolution of the 2_1^+ state energies and the $B(E2)$-values along isotopic chains $^{126-130}$Pd, $^{124-132}$Cd, $^{124-134}$Sn, $^{128-136}$Te, $^{134-138}$Xe was investigated. The calculated energies and $B(E2)$-values correctly reproduce the experimental isotopic and isotonic dependences. The structure of 2_1^+ states in the $^{126-130}$Pd and $^{124-132}$Cd nuclides were predicted. Moreover, we studied the properties of the lowest isovector collective quadrupole states. It was shown that the 2_4^+ state in ^{130}Te is the best candidate for the mixed-symmetry state.

As the second example, the isoscalar giant quadrupole resonances (ISQR) in ^{132}Sn and ^{208}Pb were studied [5]. The energy centroids and widths of the ISQR were calculated consistently with the Skyrme interaction SLy4 taking into account the coupling with large number of two-phonon configurations. This coupling results in essential increasing of the ISQR width in comparison with the QRPA calculations (see Fig. 1). We described the experimental data for ^{208}Pb and gave predictions for ^{132}Sn.

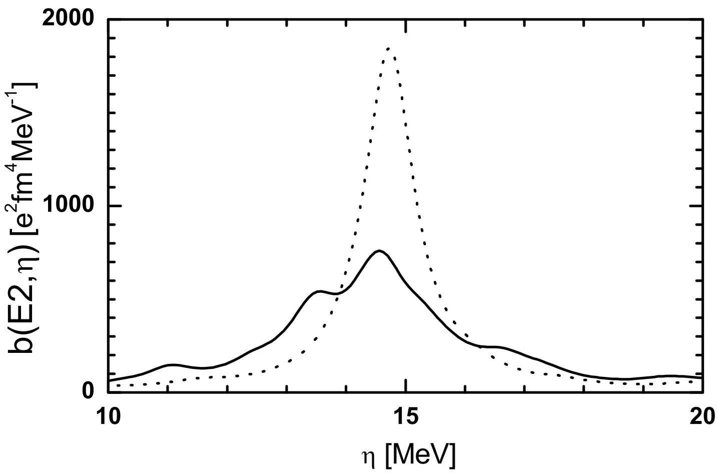

Figure 1: The quadrupole strength distribution in ^{132}Sn

In Ref. [6], two methods of elimination of the effect of the spurious state on the $E1$-transition strength distribution were compared. The first method is the "standard" one: it

Table 1: Results of calculations for Gamow-Teller excitations

	Skyrme		LM	
	E (MeV)	$B^-(GT)$	E (MeV)	$B^-(GT)$
^{48}Ca	3.95	5.75	3.03	5.82
	12.25	17.83	11.49	17.83
^{90}Zr	9.39	7.70	8.62	7.78
	17.15	22.03	16.49	21.94
^{132}Sn	3.31	1.52	3.07	1.56
	7.02	2.27	6.72	3.34
	8.18	16.25	7.96	14.97
	10.50	3.40	10.39	5.08
	15.96	69.29	14.76	67.72

uses the nucleon effective charges N/A and $-Z/A$ in the electric dipole transition operator. The second method is based on the procedure of orthogonalization of the spurious state to all one-phonon 1^- QRPA states [7]. The $B(E1)$ distributions in ^{100,124,130,132}Sn and ^{208}Pb were calculated within both the methods and very close results were obtained. The calculated positions of the $E1$ resonances agree well with the available experimental data. The self-consistent Skyrme-QRPA approach with separabelized residual interactions was also applied to the Gamow-Teller (GT) strength distributions [5]. In particular, we compared the energies and transition strengths of the charge-exchange 1^+ states calculated with the Landau-Migdal (LM) approximation for the Skyrme particle-hole interaction with those obtained with the full Skyrme force. The ^{48}Ca, ^{90}Zr and ^{132}Sn nuclides were selected as illustrative examples. As one can see from Table 1, the LM approximation properly reproduces the results of the full Skyrme interactions. Note also that the theory is in reasonable agreement with the experimental data.

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SPIN-FLIP $M1$ GIANT RESONANCE AS A CHALLENGE FOR SKYRME FORCES

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Despite a great success of the time-dependent Skyrme Hartree-Fock (TDSHF) approach in exploration of nuclear dynamics [1], it is still rarely applied to magnetic excitations, in particular to spin-flip $M1$ and scissors $M1$ giant resonances (GR). At the same time, the spin-flip $M1$ GR is an important source of knowledge on spin correlations in the Skyrme functional. The resonance also strongly depends on the spin-orbit splitting and so can serve as a robust test of the spin-orbit interaction. Besides, the spin-flip $M1$ GR is closely related to the Gamow-Teller (GT) resonance and its satisfactory treatment is relevant for the correct description of GT mode as well.

Our recent studies have shown that TDSHF has serious troubles in description of the spin-flip $M1$ GR [2, 3]. The results for different Skyrme parameterizations are contradictory and do not reproduce the experimental data. In particular, it is quite difficult to describe with one and the same Skyrme force a one-peak gross structure of $M1$ strength in doubly magic nuclei and a two-peak structure in heavy deformed nuclei. The reason of this mismatch is not the deformation splitting but instead lies in an unsatisfactory treatment of the fragile balance between spin correlations and spin-orbit interaction [2, 3].

This problem was systematically scrutinized in the framework of the self-consistent Skyrme separable random-phase-approximation (SRPA) method which, being proved as a reliable theoretical tool for investigation of electric giant resonances [4], was recently extended to magnetic excitations [2, 3]. The exploration involved 8 Skyrme forces and various (light/heavy, spherical/deformed) nuclei [2, 3]. Both isospin-mixed (for inelastic electron scattering) and isovector (T=1) channels were considered. The isovector spin-orbit interaction and tensor forces (introduced through the squared spin-orbit density J^2) were inspected as promising tools to improve the description.

It was shown that tensor forces strongly influence the spin-flip $M1$ strength and may considerably improve the simultaneous description of this GR in spherical and deformed nuclei, see Fig. 1 (bold curve). The proper choice of the Skyrme force and refit of Skyrme parameters after switching on the tensor contribution were found of crucial importance. The strong impact of tensor forces on the spin-flip $M1$ GR was also confirmed by Colò *et al.* [7]. Instead, the effect of the T=1 spin-orbit interaction (depicted in Fig. 1 for the spin-orbit parameters typical for the relativistic mean-field (RMF) model) turned out to be weak. An additional SRPA study of the spin-flip $M1$ GR (and scissors $M1$ mode) was recently performed for the chain of Nd isotopes [8]. Despite all this effort, the description of the spin-flip $M1$ GR with Skyrme forces and effect of the tensor forces still need further theoretical and experimental exploration.

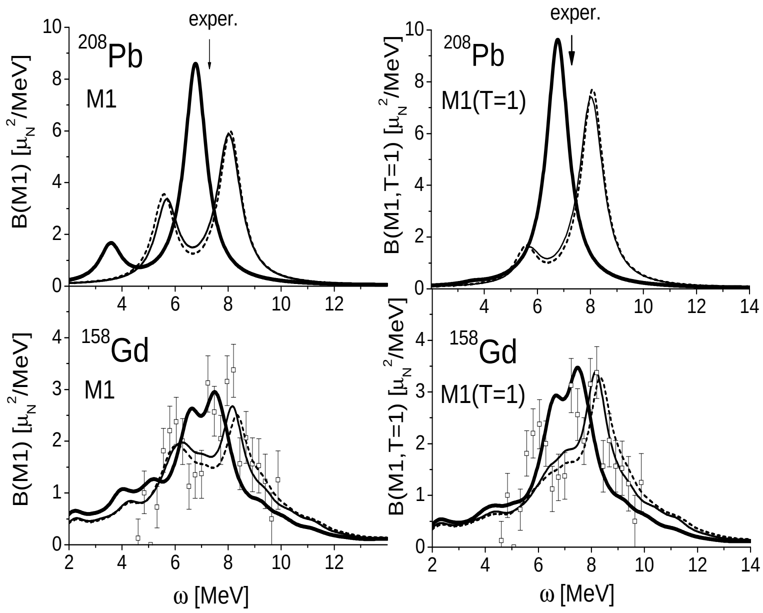

Figure 1: The isospin-mixed (left) and T=1 (right) spin-flip $M1$ strengths in ^{208}Pb and ^{158}Gd for the force SV-bas with the T=1 RMF-like spin-orbit interaction (solid curve), with the refitted tensor contribution (bold curve), and without both tensor and T=1 spin-orbit contributions (short-dash curve). The experimental data [5, 6] are exhibited by boxes with bars for ^{158}Gd and vertical arrows for ^{208}Pb. The strength is smoothed by the Lorentz weight with averaging Δ =1 MeV.

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FRAGMENTATION AND SCALES IN NUCLEAR GIANT RESONANCES

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Nuclear Giant Resonances (GR) have been the subject of numerous investigations over several decades [1]. Some of the basic features such as centroids and collectivity (in terms of the sum rules) are reasonably well understood within microscopic models. However, the question of how a collective mode like the GR dissipates its energy is one of the central issues in nuclear structure physics.

According to the accepted wisdom, GRs are essentially excited by an external field through a one-body interaction. It is natural to describe these states as collective 1p-1h states. Once excited, the GR progresses to a fully equilibrated system via direct particle emission and by coupling to more complicated configurations (2p-2h, 3p-3h, etc). The former mechanism gives rise to an escape width, while the latter yields spreading widths ($\Gamma^{\downarrow}$). An understanding of lifetime characteristics associated with the cascade of couplings and scales of fragmentations arising from this coupling remains a challenge. Recent high energy-resolution experiments of the Isoscalar Giant Quadrupole Resonance (ISQR) [2] show that the fine structure of the ISQR observed in (p, p') experiments is largely probe independent.

We propose a general approach [3] to characterize fluctuations of measured cross sections of nuclear giant resonances. Simulated cross sections are obtained from a particular, yet representative self-energy which contains all information about fragmentations. Using a wavelet analysis [4] we demonstrate the extraction of time scales of cascading decays into configurations of different complexity of the resonance. We argue that the spreading widths of collective excitations in nuclei are determined by the number of fragmentations, as seen in the power spectrum. An analytic treatment of the wavelet analysis using a Fourier expansion of the cross section confirms this principle. A simple rule for the relative life times of states associated with hierarchies of different complexity is obtained. We speculate that the fragmentations of the ISQR could be a manifestation of *self-organizing structures* [5]. Once the nuclear ISQR state is created, it is driven to an unstable hierarchy of configurations (metastable states) by quantum selection rules which connect these different complex configurations due to internal mixing. The problem of finding of these selection rules needs of course a dedicated study on its own and is left for future.

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NUCLEAR MATRIX ELEMENTS FOR NEUTRINOLESS DOUBLE BETA DECAY

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The fundamental importance of the search for the neutrinoless double beta decay ($0\nu\beta\beta$-decay) is widely accepted. After 70 years the brilliant hypothesis of Ettore Majorana is likely to be valid and is strongly supported by the discovery of neutrino oscillations and by the construction of the Grand Unified Theories. The $0\nu\beta\beta$-decay is currently the most powerful tool to test if the neutrino is a Dirac or a Majorana particle. This issue is intimately related with the origin of neutrino masses, and thus has a strong impact on astrophysics and cosmology.

The main aim of the experiments on the search for the $0\nu\beta\beta$-decay is the measurement of the effective neutrino Majorana mass $m_{\beta\beta}$. Many new projects for measurements of the $0\nu\beta\beta$-decay have been proposed with a sensitivity corresponding to $m_{\beta\beta}$ predicted under the assumption of inverted hierarchy of neutrino masses. The GERDA/MAJORANA (^{76}Ge), SuperNEMO (^{82}Se), CUORE (^{130}Te), COBRA(^{116}Cd), LUCIFER(^{82}Se), EXO(^{136}Xe), Kamland-ZEN(^{136}Xe) and other experiments hope to probe $m_{\beta\beta}$ down to 10-50 meV. These experiments would require about 1 ton of each radioactive isotope and 5-10 years of measurements.

Interpreting existing results as a measurement of $m_{\beta\beta}$ and planning new experiments depends crucially on the knowledge of the corresponding nuclear matrix elements (NMEs) that govern the decay rate. The NMEs for the $0\nu\beta\beta$-decay must be evaluated using tools of nuclear structure theory. Unfortunately, there are no observables that could be directly linked to the magnitude of the $0\nu\beta\beta$-decay nuclear matrix elements and that could be used to determine them in an essentially model independent way. The calculation of the $0\nu\beta\beta$-decay NMEs is a difficult problem because the ground and many excited states of open-shell nuclei with complicated nuclear structure have to be considered. Accurate determination of the NMEs, and a realistic estimate of their uncertainty, is of great importance. Nuclear matrix elements need to be evaluated with uncertainty of less than 30% to establish the neutrino mass spectrum and CP violating phases of the neutrino mixing matrix.

The two main approaches used for evaluation of double beta decay NMEs are the Quasiparticle Random Phase Approximation (QRPA) [1, 2] and the Large Scale Shell Model (LSSM) [3]. Both methods have the same starting point, namely, a Slater determinant of independent particles. However, there are substantial differences between both approaches, in fact the kind of correlations they include is complementary. The QRPA treats a large single particle model space, but truncates heavily the included configurations. The LSSM, by contrast, treats a small fraction of this model space, but allows the nucleons to correlate in arbitrary ways. Matrix elements for the double beta decay are calculated also by angular momentum projected (with real quasiparticle transformation) Hartree-Fock-Bogoliubov (P-HFB) wave functions [4], the Interacting Boson Model (IBM) [5] and by Energy Density Functional Method (EDF) [6]. In the P-HFB the nucleon pairs different from 0^+ in the intrinsic coordinate system are strongly suppressed compared to the results of the LSSM and the QRPA. The approaches LSSM and QRPA show also that other neutron pairs contribute strongly, which cannot be included into real P-HFB. The IBM is also restrictive: It allows only that 0^+ and 2^+ neutron pairs are changed into proton pairs.

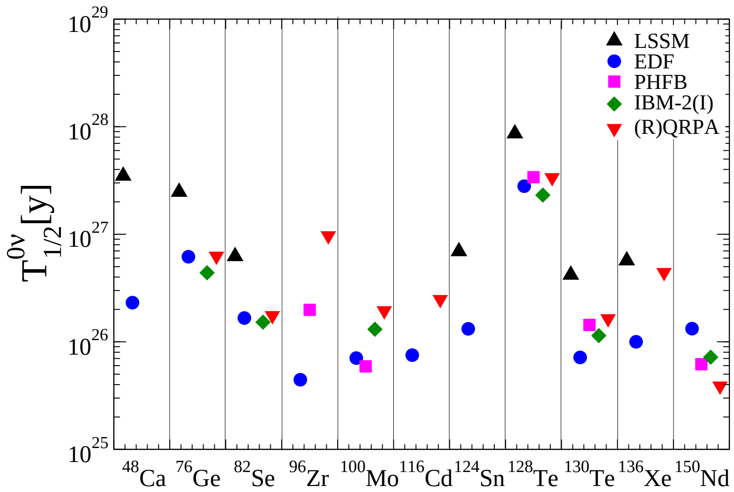

Figure 1: The calculated $0\nu\beta\beta$-decay half-lives by assuming $m_{\beta\beta} = 50$ meV and NMEs of different approaches.

Comparing the $0\nu\beta\beta$-decay nuclear matrix elements calculated using different methods gives some insight into the advantages or disadvantages of different candidate nuclei. However, matrix elements are not quite the relevant quantities. Experimentally, half-lives are measured or constrained, and the effective Majorana neutrino mass $m_{\beta\beta}$ is the ultimate goal. For $m_{\beta\beta}$ equal to 50 meV the calculated half-lives for double β-decaying nuclei of interest are presented in Fig. 1. We see that the spread of half-lives for given isotope is up to the factor of 4-5.
The improvement of the calculation of the $0\nu\beta\beta$-decay NMEs is a very important and challenging problem. The uncertainty associated with the calculation of the $0\nu\beta\beta$-decay NMEs can be diminished by suitably chosen nuclear probes. Complementary experimental information from related processes like charge-exchange reactions, muon capture and charged current (anti)neutrino-nucleus reactions is highly required. A direct confrontation of nuclear structure models with data from these processes might improve the quality of nuclear structure models [7]. The constrained parameter space of nuclear models is a promising way to reduce uncertainty in the calculated $0\nu\beta\beta$-decay NMEs [8].
The occupancies of valence neutron and proton orbits determined experimentally by J. Schiffer *et al.* represent important constraints for nuclear models used in the evaluation of the $0\nu\beta\beta$-decay NME for the ^{76}Ge $\rightarrow$ ^{76}Se transition. Clearly, having the experimental orbit occupancies available and adjusting the input to fulfill the corresponding constraint makes a difference. Within QRPA and its generalizations it was found that it was important also to choose the variant of the basic method that made such a comparison meaningful by conserving the average particle number in the correlated ground state. In [9], the conclusion was that for the ^{76}Ge $\rightarrow$ ^{76}Se transition the matrix element is smaller by 25%, reducing the previously bothersome difference with the shell model prediction

noticeably. It would be very useful to have similar constraints available also in other systems.
A microscopic state-of-the-art calculation of the NME for the $0\nu\beta\beta$-decay of ^{150}Nd with an account for nuclear deformation was performed [2, 10, 11]. The proton-neutron QRPA with a realistic residual interaction [the Brueckner G matrix derived from the charge-depending Bonn (Bonn-CD) nucleon-nucleon potential] was used as the underlying nuclear structure model. The calculated NME is suppressed by about 40% as compared with the spherical QRPA result for ^{150}Nd. By making use of this newest NME one may conclude that neutrinoless double beta decay of ^{150}Nd, to be measured soon by the SNO+ collaboration, provides one of the best probes of the Majorana neutrino mass.
Till now, Miller-Spencer Jastrow short-range correlations (SRC) have been introduced into the involved two-body transition matrix elements, changing two neutrons into two protons, to achieve healing of the correlated wave functions. In [1], the coupled cluster method (CCM) short-range correlations were considered. They were obtained as a solution of the CCM with realistic CD-Bonn and Argonne V18 interactions. By performing a consistent calculation of the $0\nu\beta\beta$-decay NMEs in which pairing, ground-state correlations and the short-range correlations originate from the same realistic NN interaction, namely, from the CD-Bonn and Argonne potentials, matrix elements for the $0\nu\beta\beta$-decay obtained are about 20% larger in magnitude when compared with the traditional approach of using the Miller-Spencer Jastrow correlations.
It is well known that there exist many mechanisms that may contribute to the $0\nu\beta\beta$-decay. By exploiting the fact that the associated nuclear matrix elements are target dependent we showed that given definite experimental results on a sufficient number of targets, one can determine or sufficiently constrain all lepton violating parameters including the mass term [12]. As a specific example we showed that assuming the observation of the $0\nu\beta\beta$-decay in three different nuclei, e.g. ^{76}Ge, ^{100}Mo and ^{130}Te, and just three lepton number violating mechanisms (light and heavy neutrino mass mechanisms as well as R-parity breaking SUSY mechanism) being active, there are only four different solutions for the lepton violating parameters, provided that they are relatively real. In particular, our analysis showed that the effective neutrino Majorana mass $|m_{\beta\beta}|$ can be almost uniquely extracted by utilizing other existing constraints (cosmological observations and tritium β-decay experiments). We also pointed out the possibility that the non-observation of the $0\nu\beta\beta$-decay for some isotopes could be in agreement with a value of $|m_{\beta\beta}|$ in sub eV region. We found that the obtained results are sensitive to accuracy of measured half-lives and to uncertainties in calculated NMEs.
In summary, there has been a significant progress in understanding the source of the spread of calculated NMEs. Nevertheless, there is no consensus as yet among nuclear theorists about their correct values and the corresponding uncertainty. However, a recent development in the field is encouraging. There is a reason to be hopeful that the uncertainty will be reduced.

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ELECTRON CAPTURE AT FINITE TEMPERATURES: FIRST-FORBIDDEN TRANSITIONS AT PRESUPERNOVA CONDITIONS

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Weak interaction mediated processes play an important role in many astrophysical scenarios. For example, electron captures (EC) on iron group nuclei initiate the gravitational collapse of the core of a massive star triggering the supernova explosion. Evaluating the rates of these processes one should take in mind that they occur at temperatures of the order of a few hundred keV to a few MeV, and thus one should consider not only transitions between the ground and excited nuclear states, but also between different excited states in different nuclei.

To date, the most reliable calculations of EC rates in stellar environment for hot nuclei from the iron region have been performed in the framework of the large-scale shell-model (LSSM) approach (see, e.g. [1] and references therein). However, the LSSM approach currently is not feasible for nuclei with neutron numbers $N > 40$ and proton numbers $20 < Z < 40$ and/or energies of captured electrons $E_e \geq 15$ MeV because of extremely large model spaces.

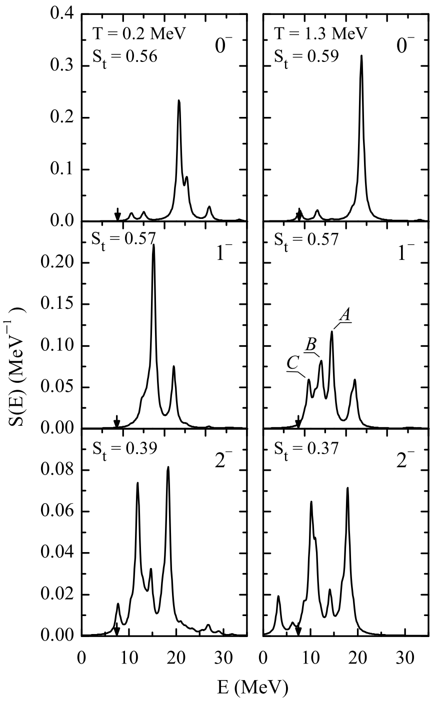

Figure 1: Folded strength distributions of $p \to n$ transitions with $\Delta J = 0^-,\ 1^-,\ 2^-$ in ^{76}Ge at $T = 0.2$ MeV (left panels) and $T = 1.3$ MeV (right panels); E is the transition energy. The strength distributions for the 2^- multipole correspond to 25 MeV electrons. S_t is the total strength. The arrows indicate the zero temperature threshold. The letters label the 1^- transitions: $A \equiv 1f^p_{7/2} \to 2d^n_{5/2}$, $B \equiv 1f^p_{5/2} \to 1g^n_{7/2}$, $C \equiv 1f^p_{7/2} \to 1g^n_{9/2}$.

During gravitational collapse the nuclear composition moves towards a higher mass number and more neutron-rich nuclei. Moreover, the Gamow-Teller transitions determine EC rates at the early stage of the collapse. At core densities $\rho > 10^{11}$ g/cm^3, when electron

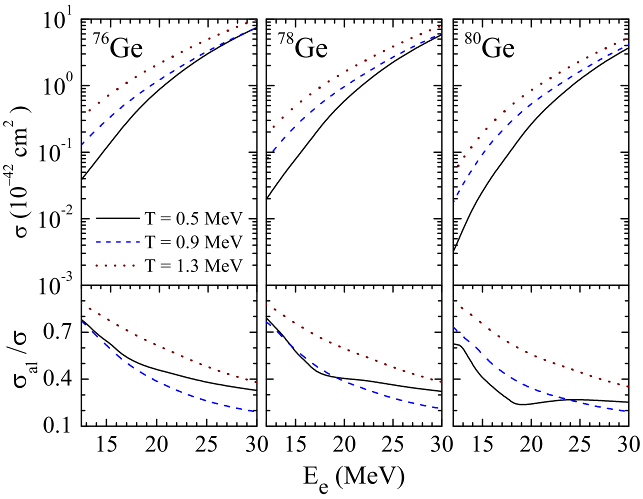

Figure 2: Upper panels: Electron capture cross sections for ^{76,78,80}Ge as functions of electron energy E_e calculated for various temperatures. Lower panels: Relative contributions of allowed transitions to the electron capture cross sections.

chemical potential reaches the value ∼18 MeV, the forbidden EC transitions can no longer be neglected.

In Refs. [2, 3], we considered the first-forbidden EC on the hot Ge isotopes within the approach combining the thermal QRPA in the context of the thermo field dynamics [4, 5] and the Qusiparticle-Phonon Model.

The charge-exchange spin-dipole strength distributions were calculated with the Hamiltonian of the Qusiparticle-Phonon Model consisting of the phenomenological single-particle Wood-Saxon potential, the BCS pairing interaction and separable dipole and spin-dipole forces [5].

The strength distributions of charge exchange first-forbidden transitions with $\Delta J = 0^-$, 1^-, 2^- in ^{76}Ge are shown in Fig. 1 for two temperatures T. As it is seen from the figure, a temperature increase weakly affects the main peaks of the 0^-, 2^- strength distributions but induces a significant spread in the 1^- strength distribution. This effect is a result of intimate interference between particle-hole and particle-particle, hole-hole 1^- configurations caused by thermal smearing of the Fermi surface [3].

In Fig. 2, the EC cross sections for ^{76,78,80}Ge are shown for different T. The temperature dependence of the cross sections is most pronounced at moderate electron energies $E_e \leq 15$ MeV where the Gamow-Teller transitions dominate. For larger electron energies, the first forbidden transitions become increasingly important. As the strength of the first forbidden transitions is less sensitive to temperature change, the capture cross sections at $E_e \sim 30$ MeV depend only weakly on temperature.

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FORMATION OF HYPERDEFORMED STATES IN THE ENTRANCE CHANNEL OF HEAVY-ION REACTIONS

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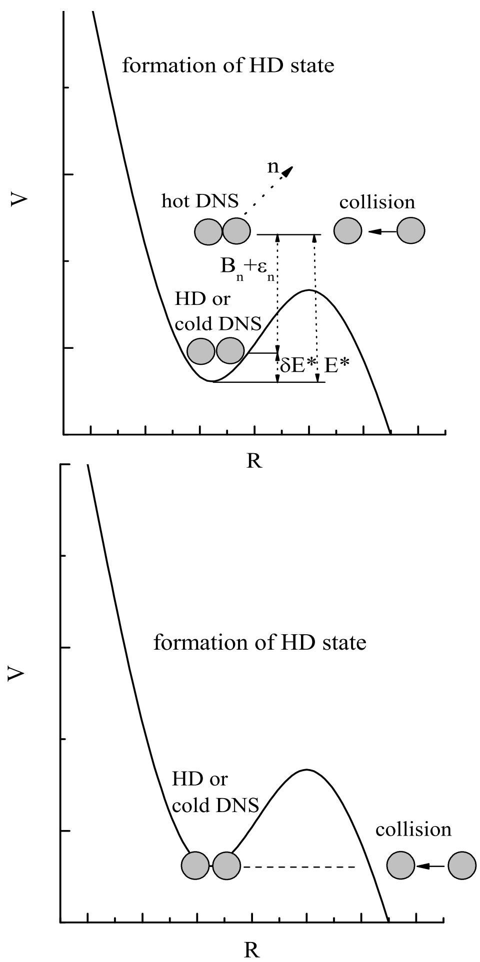
Hyperdeformed (HD) states are highly elongated exotic nuclear shapes which are caused by the third minimum in the potential energy surfaces (PES) which appears at very large quadrupole deformation parameters $\beta_2 \gtrsim 0.9$. The evidence of low-spin HD-states in actinides has been experimentally established in induced fission reactions (n,f), (t,pf), and (d,pf). The question of experimental indications of high-spin HD-state is still open. According to the cluster interpretation, HD state can be considered as a dinuclear system (DNS) of two clusters in a touching configuration. The relative distance between the centers of the clusters corresponds to the minimum of the nucleus-nucleus interaction potential. The minimum of nucleus-nucleus potential energy contains the quasibound states with the energies below the potential barrier and with large half-lives.



Using the cluster approach we proposed a model of the HD state formation in the entrance channel of heavy-ion reaction at bombarding energies near and below the Coulomb barrier. The initial excited DNS can then be de-excited by the emission of a neutron to the cold quasibound state which is identical to the HD state. Another mechanism for the population of the HD state is the direct sub-barrier tunneling (see the Figure). In this kind of reactions the high-spin HD states can be populated and experimentally identified. The neutron emission from the initial excited DNS, which competes with the quasifission and the diffusion of the initial DNS to more symmetric or asymmetric configurations, is described by using a statistical approach. Tunneling through the Coulomb barrier is considered using the quantum diffusion approach with the formalism of reduced density matrix. The experimental identification of the HD state can be obtained by measuring the consecutive col-

lective rotational $E2-$transitions in the HD band in coincidence with the decay fragments of the DNS trapped in the HD minimum.
The optimal reactions and conditions (bombarding energies, range of angular momenta) for the identification of the HD states are proposed, and the HD state formation and identification cross sections are estimated. At bombarding energies near the Coulomb barrier we propose to consider the reactions ^{48}Ca+^{124,128,130,132,134}Sn, ^{48}Ca+^{136,138}Xe, ^{48}Ca+^{137,138,140}Ba, ^{40}Ca+^{83,84}Kr, ^{48}Ca+^{83,84,86}Kr, ^{40,48}Ca+^{40,48}Ca, ^{58,60}Ni+^{58,60}Ni, and ^{40}Ca+^{58}Ni as good candidates for the production and experimental identification of the HD states. The estimated identification cross sections for the HD states formed in these reactions are of the order of 1 nb – 2.5 μb for optimal bombarding energies and range of angular momenta. We propose to consider the reactions ^{48}Ca+^{124}Sn, ^{48}Ca+^{136}Xe, ^{48}Ca+^{138}Ba, ^{48}Ca+^{140}Ce, ^{48}Ca+^{86}Kr, ^{58}Ni+^{58}Ni, ^{40}Ca+^{40}Ca, and ^{48}Ca+^{48}Ca as good candidates for the production and experimental identification of the HD states at the sub-barrier energies. The estimated maximal values of the partial HD identification cross sections for these reactions vary from 0.1 nb up to 0.5 mb.

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PRODUCTION OF EXOTIC NUCLEI IN TRANSFER-TYPE REACTIONS

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Besides the reactions at intermediate energies the multinucleon transfer and quasifission-type reactions at low energies are actively discussed to produce exotic nuclei. As it was showed, the diffusive multinucleon transfer-type reactions can be described as an evolution of the dinuclear system (DNS) which is formed in the entrance channel of the reaction after dissipation of the kinetic energy and angular momentum of the relative motion. The dynamics of the process is considered as a diffusion of the DNS in the charge and mass asymmetry coordinates which are defined here by the charge and neutron numbers Z and N of the light nucleus of the DNS. During the evolution in the charge and mass asymmetry coordinates the excited DNS can decay into two fragments at a relative distance R between the centers of the DNS nuclei. So within the DNS model the production of the exotic nucleus is treated as a three-step process. First, the initial DNS with the light nucleus $(Z_i,\ N_i)$ is formed in the peripheral collision for a short time. Second, the DNS with the light exotic nucleus (Z, N) is produced by nucleon transfers. Then this DNS separates into two fragments.

The suggested method is suitable to predict the mass and charge yields and the production cross sections for certain products of multinucleon transfer reactions. The calculated production cross sections of the neutron-rich isotopes in the reactions ^{48}Ca+^{238}U,^{244}Pu at incident energies near the Coulomb barrier are presented in Figs. 1 and 2. We treat only the reactions leading to excitation energies of light neutron-rich nuclei equal to or smaller than their neutron separation energies ($E^*_L(Z,N,J) \leq S_n(Z,N)$). In this case, W_{sur}=1 and the primary and secondary yields coincide. In Figs. 1 and 2, the values of $E_{\mathrm{c.m.}}$ provide the condition $E^*_L(Z,N,J) = S_n(Z,N)$. The predicted values of $S_n(Z,N)$ for unknown nuclei are taken from the finite range liquid drop model. If $E^*_L(Z,N,J) > S_n(Z,N)$, the primary neutron-rich nuclei are transformed into the secondary nuclei with a smaller number of neutrons because of the de-excitation by neutron emission. The DNS evolution in the reactions treated can be schematically presented in the following way: ^{48}Ca+^{238}U$\rightarrow$^{78,80}Zn+^{208,206}Pb$\rightarrow$^{82,84,86}Zn+^{204,202,200}Pb and ^{48}Ca+^{244}Pu$\rightarrow$^{84,82}Ge+^{208,210}Pb$\rightarrow$^{86,88,90,92}Ge+^{206,204,202,200}Pb. The system initially moves to the deep minimum of the potential energy surface (energetically favorable) which is caused by the shell effects around the DNS with the magic heavy ^{208}Pb and light ^{80}Zn or ^{82}Ge nuclei; then from this minimum it reaches the DNS with the exotic light nucleus by fluctuations in mass asymmetry. For low excitation energy, the evolution of the dinuclear system towards symmetry is hindered by this minimum.

The production cross sections of the primary isotopes in the reactions ^{48}Ca+^{238}U at $E_{\mathrm{c.m.}}$=189 MeV were calculated. The primary neutron-rich nuclei of interest are excited and transformed into the secondary nuclei with a smaller number of neutrons without a loss of the cross section because the neutron emission is dominant over other deexcitation channels. Since the predicted production cross sections for the new exotic isotopes ^{193}W, ^{195,196}Re, ^{198}Os, and ^{200}Ir are at the microbarn level, they can be easily identified. For these nuclei, the known heaviest isotopes are in the vicinities of maxima of the primary isotopic distributions. Since the calculated production cross sections for the new

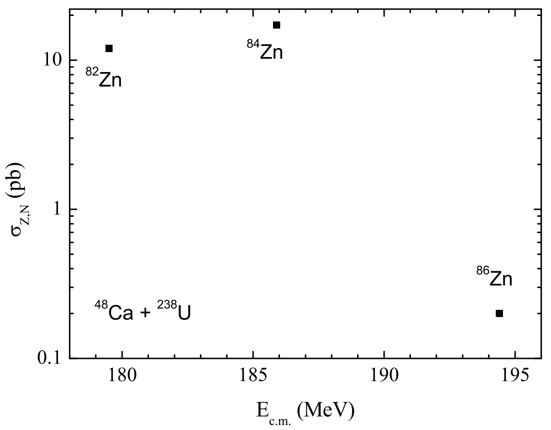

Figure 1: The expected cross sections for the indicated neutron-rich isotopes of Zn produced in the ^{48}Ca+^{238}U reaction at values of $E_{\text{c.m.}}$ providing the excitations of these isotopes to be equal to the corresponding thresholds for the neutron emission.

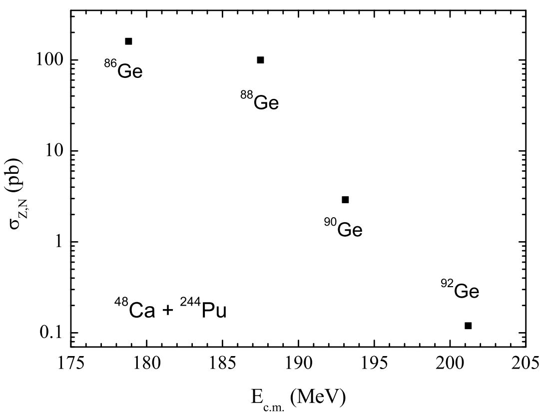

Figure 2: The same as in Fig. 1, but for the indicated neutron-rich isotopes of Ge produced in the ^{48}Ca+^{244}Pu reaction.

exotic isotopes ^{178}Er, ^{180,181}Tm, $^{182-184}$Yb, $^{185-187}$Lu, ^{190}Hf, $^{191-193}$Ta, ^{194,196}W, ^{197,199}Re, ^{199,200}Os, ^{201,202}Ir, ^{203}Pt are between the microbarn and nanobarn levels, they can also be detected with the present experimental setups.

In the quasifission reactions ^{48}Ca+^{244,246,248}Cm at beam energies close to the corresponding Coulomb barriers, one can produce the new isotopes of superheavies with $Z = 103 - 108$, which undergo fission (the fission width is much larger than the neutron emission width). The calculated results indicate that these quasifission reactions provide a very efficient tool for the study of new isotopes of superheavy nuclei that fill the gap between the isotopes produced in the cold and hot complete fusion reactions. The predicted cross sections of the fission, which follow multinucleon transfer are at the level (100 nb-100 μb). One can propose the experiments on the quasiternary fission in which the fission fragment mass and the angular distributions in coincidence with the complementary transfer products can be measured.

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INVESTIGATION OF HINDRANCE TO FUSION TO SELECT REACTIONS FOR SYNTHESIS OF SUPERHEAVY ELEMENTS

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The observed evaporation residues in experiments are a result of the de-excitation of a heated and rotating compound nucleus formed in competition of a complete fusion with quasifission and fast fission processes. The last two processes are the hindrance to formation of the compound nucleus which is the necessary condition to observe evaporation residues being registered as a superheavy element with the total charge of reacting nuclei larger than 110 [1, 2]. The correct estimation of the fusion cross section in the reactions with massive nuclei is a difficult task from both theoretical and experimental points of view.

Different assumptions about the fusion process are used in different theoretical models and they can predict different cross sections [3-5]. The experimental methods used to estimate the fusion probability depend on the unambiguity of identification of the complete fusion reaction products among the quasifission products. The difficulties arise when the mass (charge) and angular distributions of the quasifission and fusion-fission fragments strongly overlap depending on the reaction dynamics [6, 7]. The comparison of our results obtained in the framework of dinuclear system (DNS) model with the experimental data for the ^{48}Ca+^{154}Sm reaction [6] showed (see the left panel of Fig. 1) that the yield of measured fission-like fragments (stars) at the large bombarding energies was higher than the theoretical fusion-fission cross section (the dash-double-dotted line). This deviation is explained by mixing the contributions of the quasifission (the short-dashed line) and the fast fission (the dash-dotted line) fragments into the measured data of fission-like products [8]. Therefore, the experimental quasifission cross sections (triangles) are lower than the theoretical ones (the short-dash line). The small value of the contribution of the theoretical fusion-fission at low energies is due to the high fission barrier (B_{f}=12.33 MeV) for the compound ^{202}Pb nucleus.

The other consequence of the unintentional inclusion of the quasifission and fast fission contributions in the fission-like fragment yields for a correct estimation of the fusion cross section is demonstrated in the right panel of Fig. 1. Analyzing the ^{19}F+^{181}Ta and ^{16}O+^{184}W reactions the authors of Ref. [9] have concluded from the comparison of the evaporation residue cross sections (normalized to the fusion cross sections) that the difference between the corresponding data at high excitation energies is due to the difference in the maximal values of angular momentum distributions. Our theoretical estimations showed that the maximum values of angular momentum distributions in these reactions are nearly the same and the difference is caused by the increase of the quasifission contributions in the ^{19}F+^{181}Ta reaction into the reconstructed experimental fusion cross section in the normalizing procedure (see the right ordinate axis in the right panel of Fig. 1). So the difficulties in selecting the quasifission and fusion-fission products from the measured data can lead to an incorrect interpretation of the physical results. The importance of knowledge of the realistic fusion cross sections was demonstrated at our suggestion to employ the ^{54}Cr+^{248}Cm reaction for production of the Z=120 element since it is the most favorable one in comparison with the ^{58}Fe+^{244}Pu and ^{64}Ni+^{238}U reactions [8].

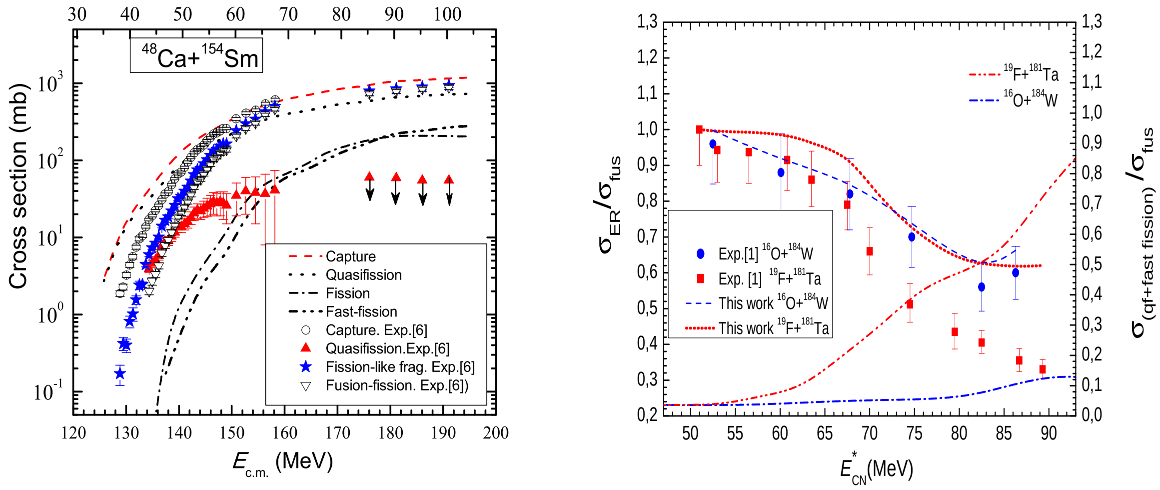

Figure 1: (Left panel: Comparison of our results obtained in the framework of the DNS model for the capture, complete fusion, quasifission, fast-fission and evaporation residue cross sections with the data of the fusion-fission and quasifission from Ref. [6]. Right panel: Comparison of the experimental evaporation residue cross sections (normalized to the fusion cross sections) for the systems ^{16}O+^{184}W (solid circles) and ^{19}F+^{181}Ta (solid squares) [9] with the corresponding theoretical results (dashed and dotted lines, respectively) depending on the excitation energy $E^*_{\rm CN}$ of a compound nucleus (CN) (left axis). The theoretical sum of the quasifission and fast fission cross sections (normalized to the fusion cross sections) for the ^{16}O+^{184}W (dot-dashed line) and ^{19}F+^{181}Ta (dot-dot-dashed line) systems is presented versus $E^*_{\rm CN}$ and compared on the right axis.

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HALO FORMATION AND BREAKUP

S. N. Ershov

The evolution of nuclear structure from the valley of stability to the limits of nuclear existence (driplines) and beyond, is one of the most important and interesting topics in modern nuclear physics. Remarkable phenomena are observed in nuclei near driplines including a new type of nuclear structure, called halos, identified in neutron-rich weakly bound light nuclei. Characteristic features of halo systems are extreme few-body clusterization and extraordinary large sizes. Two-neutron halo nuclei, like ^{6}He, ^{11}Li and ^{14}Be, display the most widespread exotic type of halo phenomenon: They are Borromean, meaning that they decay into three constituent fragments when excited above the lowest threshold. Studies of correlations in relative motions between the three fragments, open a way for extended exploration of halo structure, its formation and how it dissolves. This demands a clear understanding of both nuclear structure and the reaction mechanism, inducing the breakup.
Fragment correlations are accessible via different cross sections that can be measured if fragments are detected in coincidence. Different correlations contain different information about nuclear structure and reaction dynamics. In general, continuum halo excitations at different excitation energies are coupled by reaction dynamics. However, there are physical situations when breakup via a one-step excitation mechanism is most favorable and also simple enough to allow theoretical modeling. The microscopic four-body distorted wave theory for two-neutron halo breakup reactions leading to low-lying halo excitations was developed [1, 2], which accounts for both elastic and inelastic breakup. The Coulomb and nuclear dissociation is included in a consistent way. The method of hyperspherical harmonics is used for a consistent description of specific features of the halo structure of the ground state and the fragment motion in the continuum.

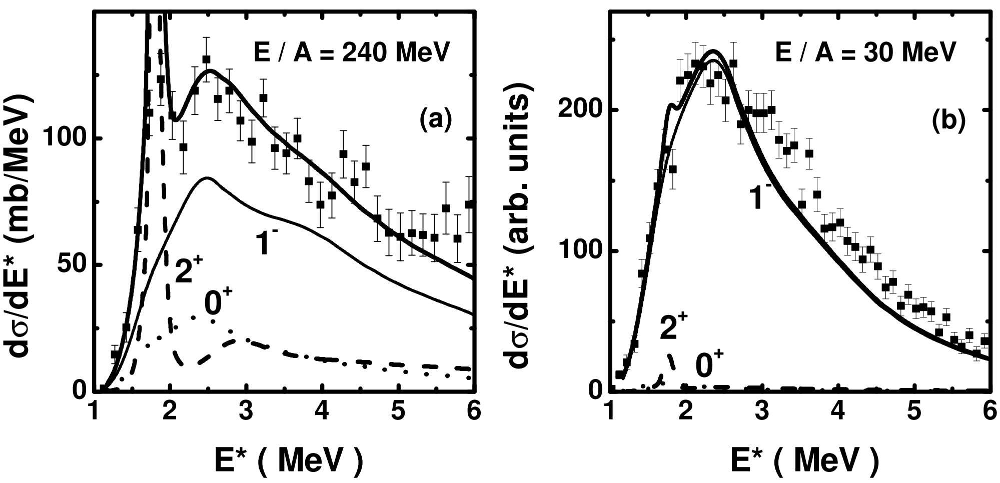

Figure 1: Comparison of theoretical ^{6}He excitation spectra (thick solid curves) for the ^{6}He + ^{208}Pb *Coulomb dominated* breakup at high and low collision energies with experimental data, and the dipole 1^-, quadrupole 2^+, and monopole 0^+ contributions.

Description of the excitation spectrum of a cluster nucleus in few-body breakup is a key objective for continuum spectroscopy. A practical way to test theoretical assumptions is to calculate the spectrum using the same structure model for different reaction mechanisms. Figure 1 shows the comparison of the calculated ^{6}He excitation spectrum (thick solid curve) with experimental data for breakup reactions on lead target at collision energies 240 and 30 MeV/nucleon obtained at GSI [3] and GANIL [4] respectively. Theory reproduces the shape of the low-lying excitation spectrum for both collision energies, though the breakup mechanisms are quite different. Since the target is a heavy nucleus, Coulomb dissociation dominates at both energies. At low collision energy the dipole mode dominates the spectrum and only a small remnant of the 2^+ (1.8 MeV) three-body resonance is visible. At high collision energy the dipole mode remains large but quadrupole and monopole excitations give considerable contributions to the cross section, in particular the quadrupole resonance.The calculations describe the experimental data for fragment correlations near the breakup threshold rather well. Experimental data are called for exclusive cross sections since theory now provides correlation cross sections from fully inclusive to fully exclusive (spectrum).

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STUDY OF ^{6}He + ^{12}C ELASTIC SCATTERING USING A MICROSCOPIC OPTICAL POTENTIAL

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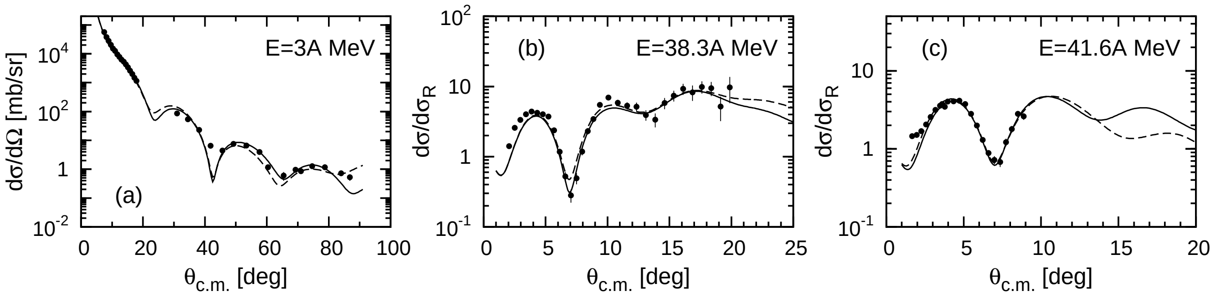
The ^{6}He+^{12}C elastic scattering data at beam energies of 3, 38.3 and 41.6 MeV/nucleon were analyzed theoretically by utilizing the microscopic optical potentials (OP) [1,2]. Optical potentials can be selected with different forms of a surface term, and the mostly appropriate one was established as follows:

$$U_{opt} \;=\; N_R V^{DF}(r) \,+\, iN_I W(r) \,-\, iN_I^{sf} r^2 (dW/dr). \tag{1}$$

Here the real part $V^{DF}(r)$ is the standard double-folding potential [3] while for the imaginary part the usage was made of both the shapes — $W = V^{DF}(r)$ and that inherent in the high-energy approximation $W = W^H(r)$ [4]. Calculations of OP's are based on the unfolded neutron and proton density distributions for ^{6}He from [5] and for ^{12}C from [6]. The problem of ambiguity of the obtained set of OPs was resolved by selection of only those which obey the known energy dependence of the respective volume integrals $ReJ(E) + iImJ(E) \,=\, -(4\pi/A_pA_t)\int U_{opt}(r)r^2dr$. In the Figure the calculated elastic scattering cross sections of ^{6}He+^{12}C are shown where solid curves are for $W = W^H$ and dash-dotted for $W = V^{DF}$. Conclusions read that (a) the foregoing OP explains the ex-



perimental data fairly well fitting only two or three parameters $\{N\}$ having the meaning of a strength of OP; (b) the role of the OP surface term is revealed only at higher energies; (c) special attention should be paid to the Pauli blocking effect and contributions of breakup channels to ImU$_{opt}$ while the exchange effects have already been accounted for in the real part of OP [3].

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PECULIARITIES OF THE THREE-BODY WAVE FUNCTIONS NEAR THE TRIPLE IMPACT POINT

V. V. Pupyshev

The main aim of this note is to present our most interesting results described in detail in Section 2 of the recent review [1]. To this end, we remind the basic definitions and formulae.

Let $\{p_1, p_2, p_3\}$ by a system of three quantum particles p_1, p_2 and p_3.

For this system, as the relative coordinates, we use the Jacobi vectors $(\mathbf{x}_i, \mathbf{y}_i)$, $i = 1, 2, 3$, and the corresponding sets of the hyperspherical coordinates (r, Ω_i): the hyperradius $r = (\, x_i^2 + y_i^2 \,)^{1/2}$ and the hyperangles $\Omega_i = (\, \hat{x}_i, \hat{y}_i, \varphi_i \,)$, where $\varphi_i \equiv \mathrm{atan}(y_i/x_i)$.

The point $(r = 0, \forall\Omega_i)$ is called the triple impact point.

By assumption, all interactions in the system $\{p_1, p_2, p_3\}$ are the two-body and central potentials of a wider class than the Coulomb potentials:

$$V_i(x_i) = q_i/x_i + \tilde{V}_i(x_i) = \sum_{n=-1}^{\infty} V_{in}\, x_i^n, \quad x_i \to 0\,; \quad q_i, V_{in} = \mathrm{const}; \quad i = 1, 2, 3.$$

For these interactions the total set ε of the conserved quantum numbers contains the total angular momentum ℓ, the magnetic number m and the total spacial parity σ. Let Ψ^ε be the wave function of the system $\{p_1, p_2, p_3\}$ having the set ε.

In the Faddeev theory this wave-function is represented as

$$\Psi^\varepsilon(\mathbf{x}_i, \mathbf{y}_i) = \Psi_i^\varepsilon(\mathbf{x}_i, \mathbf{y}_i) + \sum_{k \neq i} \Psi_k^\varepsilon(\mathbf{x}_k(\mathbf{x}_i, \mathbf{y}_i), \mathbf{y}_k(\mathbf{x}_i, \mathbf{y}_i))$$

and the components Ψ_i^ε satisfy the system of differential equations

$$[\, H_0(\mathbf{x}_i, \mathbf{y}_i) - E \,]\ \Psi_i^\varepsilon(\mathbf{x}_i, \mathbf{y}_i) = -V_i(x_i)\, \Psi^\varepsilon(\mathbf{x}_i, \mathbf{y}_i)\,,$$

where H_0 and E are the free Hamiltonian and the total energy.

Using the expansions

$$\Psi_i^\varepsilon(r, \Omega_i) = \frac{1}{2} \sum_{a,b} \left[\, \sigma + (-1)^{a+b} \,\right]\ \Psi_{iab}(r, \varphi_i)\, \mathcal{Y}_{ab}^{\ell m}(\hat{x}_i, \hat{y}_i)$$

over the bispherical harmonics $\mathcal{Y}_{ab}^{\ell m}(\hat{x}_i, \hat{y}_i)$ we proved the representations for the searched components Ψ_{iab}^ε in the form of the Fock-type series

$$\Psi_{iab}(r, \varphi_i) = \sum_{n=0}^{\infty} r^n \sum_{m=0}^{M(n)} s^m\, \Phi_{iab}^{nm}(\varphi_i)\,, \quad s \equiv \ln r; \quad i = 1, 2, 3.$$

Then we showed that the angular functions $\Phi_{iab}^{nm}(\varphi_i)$ are uniquely defined by the recurrence chain of the second order differential equations with the homogeneous boundary conditions. Finally, we constructed the Fock-type series for the wave function Ψ^ε as the sum of the above-mentioned Fock-type series for three components Ψ_i^ε.

Now one can discuss three peculiarities which we have found for the first time.

The first peculiarity is the dependence of the limit $M(n)$ of the Fock-type series on the structure of the two-body interactions: $M(n) = [n/2]$, when $q_i \neq 0$, if $q_i = 0$ but $V_{i1} \neq 0$, then $M(n) = [n/6]$, in the case $V_{in} = 0$, $n = 0, 1, 3, \ldots$, the limit $M(n) = 0$ for any n.
The second peculiarity is the dependence of the wave function Ψ^{ε} on the total angular momentum ℓ and the total spacial parity σ. To illustrate this peculiarity, we present explicit asymptotics of this function as $r \to 0$ in two cases.
In both cases the symbols f and X, B stands for the known function and the numerical coefficients, Y_{Lab} are the three-body hyperharmonics, and, finally, $L \equiv \ell$ for normal parity $\sigma = (-1)^{\ell}$ and $L \equiv \ell + 1$ for abnormal parity $\sigma = (-1)^{\ell+1}$.
In the case $q_i \neq 0$ the wave function has the asymptotics containing the term $O(r^2 s)$:

$$\begin{aligned}\Psi^{\varepsilon}(r, \Omega_i) &= r^L \sum_{a+b=L} \big[\, X^L_{ab}\, Y^{\ell m}_{Lab}(\Omega_i) + 2\, r \operatorname{cosec} 2\varphi_i\, f^{L1}_{ab}(\varphi_i)\, \mathcal{Y}^{\ell m}_{ab}(\hat{x}_i, \hat{y}_i) + \\ &+ \; r^2 s\, B^{L2}_{ab}\, Y_{L+2,ab}(\Omega_i) + O(r^2) \,\big] \,.\end{aligned}$$

In the case $q_i = 0$, $V_{i1} \neq 0$ the asymptotics of the wave function does not contain this term and reads as

$$\begin{aligned}\Psi^{\varepsilon}(r, \Omega_i) &= r^L \sum_{a+b=L} \big\{ \big[\, X^L_{ab} + r^2\, F^{L2}_{Lab} \,\big]\, Y^{\ell m}_{Lab}(\Omega_i) + 2\, r^3 \operatorname{cosec} 2\varphi\, f^{L3}_{ab}(\varphi_i)\, \mathcal{Y}^{\ell m}_{ab}(\hat{x}_i, \hat{y}_i) \big\} + \\ &+ \; r^{L+2} \sum_{a+b=L}^{L+2} X^{L+2}_{ab}\, Y^{\ell m}_{L+2,ab}(\Omega_i) + O(r^{L+4}) \,.\end{aligned}$$

The third peculiarity means that the asymptotics of the wave function has the different functional form in the cone-type regions, whose boundaries are defined by the values of the kinematical angles γ_{ki}.
To clarify this peculiarity, we analyze the simplest case when $q_i \neq 0$ and $\ell, m = 0$; $\sigma = 1$. In this case, the wave function has the asymptotics

$$\Psi^{\varepsilon}(x_i, y_i) = \left\{ X \left[2 + x_i\, q_i + \sum_{k \neq i} q_k\, g(x_i, y_i; \gamma_{ki}) \right] + 4B\, (x_i^2 - y_i^2)\, s \right\} + O(r^2)\,,$$

in which the function g is expressed in terms of the functions $\tilde{s} \equiv |\sin \gamma_{ki}|$ and $c \equiv \cos \gamma_{ki}$

$$g(x_i, y_i; \gamma_{ki}) \equiv \begin{cases} c x_i + (\tilde{s} y_i)^2/(3 c x_i)\,, & y_i \leq x_i \operatorname{ctg}\gamma_{ki}\,; \\ \tilde{s} y_i + (c x_i)^2/(3 \tilde{s} y_i)\,, & y_i \geq x_i \operatorname{ctg}\gamma_{ki}\,. \end{cases}$$

As one sees, there exist two particular rays $\varphi_i = \gamma_{ki}$, $k \neq i$ through which the passage changes the functional form of the function g. Hence, the asymptotics of the wave function Ψ^{ε} has the same property.
As we have proven, the structure of the three-body wave function near the triple impact point depends on the structure of the pair-interactions, the values of the total angular momentum and parity, and the subregions of the configuration space whose boundaries are defined only by the values of the kinematical angles.

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LOW-DIMENSIONAL FEW-BODY PHYSICS OF ULTRACOLD ATOMS AND MOLECULES

V. S. Melezhik

Low-dimensional quantum systems have recently become experimentally accessible with impressive development of the physics of ultracold atoms and molecules [1, 2]. It has stimulated the necessity of more detailed and deep investigations in low-dimensional few-body physics. Different aspects here demand investigation and become actual. Thus, free-space scattering theory is no longer valid and the development of quantum scattering theory in low-dimensions, including effects of confining geometry, is needed [3].
During the 2009 and 2010 years we found and investigated two novel effects [4, 5] in the ultracold atomic collisions in harmonic traps. These investigations were performed in collaborative work with theoreticians from Hamburg University (group of P. Schmelcher) and experimentalists from Innsbruck University (group of H.C. Nägerl).
We analyzed the quantum dynamics of heteronuclear atomic collisions in harmonic waveguides and suggested a novel mechanism for the resonant formation of polar molecules [4]. We showed that molecular formation rates can be tuned by changing the trap frequencies ω_1 and ω_2 characterizing the transverse modes of the atomic species. The origin of this effect is the confinement-induced mixing

$$\Delta V(\rho_{CM}, \mathbf{r}) = \mu(\omega_1^2 - \omega_2^2)\rho_{CM} r \sin\theta \cos\phi \tag{1}$$

of the relative $\mathbf{r} = \{r, \theta, \phi\}$ and center-of-mass ρ_{CM} variables in the atomic collision leading to a coupling of the diatomic continuum to CM excited molecular states in closed transverse channels.

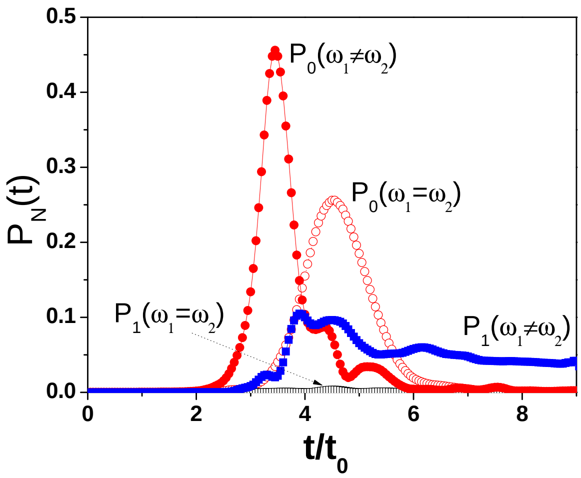

Figure 1: Illustration of the resonant molecule formation. Calculated evolution in time of the population probabilities $\mathrm{P}_N(t)$ of the final molecular states without ($N = 0$) and with ($N = 1$) the CM excitation. It is shown that in the case $\omega_1 = \omega_2$ of collision of the identical atoms the coupling (1) is absent and the effect of molecular formation in the excited CM state ($N = 1$) vanishes. Time is given in units of $t_0 = \pi/\omega_2$.
The confinement-induced resonances (CIRs) were observed in strongly interacting quantum-gas systems with tunable interactions for 1D- and 2D-geometry of confining optical potentials [5]. In the 1D-system with transverse confinement CIRs are caused by a coupling between the incident channel of two colliding atoms and the closed channel with a transversally excited molecular state (see Fig.2(a)). It was observed by characteristic atomic loss

and heating signatures that atom-atom scattering was modified substantially under the condition of the CIR appearance when the s-wave scattering length a_{3D} approached the length scale $a_\perp$ associated with the confining potential. The prediction by V.Melezhik was also confirmed that introducing an anisotropy in the transversal confinement ($\omega_1 \neq \omega_2$) leads to the CIR splitting (see Fig.2(d)). The effect is a consequence of lifting the degeneracy of the threshold of the closed channel with a transversally excited molecular state if $\omega_1 \neq \omega_2$ (see Fig.2(b)). The appearance of additional resonances was observed with increasing anisotropy. In the limit of a 2D system (the case of very strong anisotropy $\omega_1 \gg \omega_2$) one resonance was found to survive.

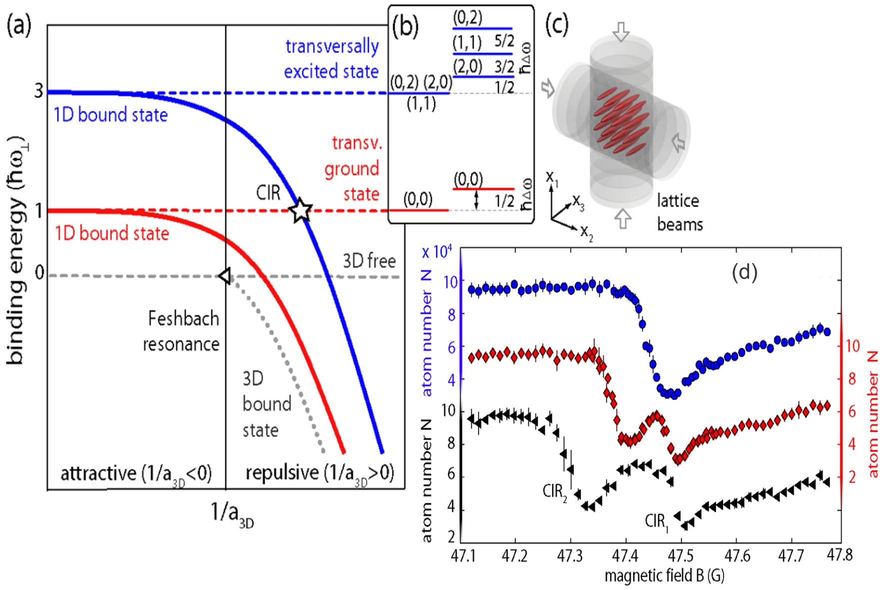

Figure 2: Illustration of the mechanism responsible for a CIR (a-b), the experimental setup (c) and the result of experiment on detection of CIRs (d). (a) The energy levels near the scattering resonance are plotted as a function of $1/a_{3D}$. The CIR occurs for $a_{3D} \simeq a_\perp$ when scattering atoms are allowed to couple to transversally excited bound states [3]. (b) indicates the shift and splitting for anisotropic confinement characterized by $\Delta\omega = \omega_2 - \omega_1$. (c) Two laser beams create an optical lattice that confines the atoms to an array of approximately 3000 independent, horizontally-oriented elongated 1D tubes. (d) Splitting of the CIR for transversally anisotropic confinement at ω_2/ω_1 =1.1 and 1.18.

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NEW MESON-NUCLEUS FEW BODY SYSTEMS

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Studies of interactions of ϕ-mesons with nuclear systems are interesting for the following reasons:

1. Formation of a new nuclear cluster.

   As it was shown in the different models of elementary ϕN interaction, there is rather strong attraction between them at low energies. From this point of view it seems to be interesting to study a possibility of existence of 3-particle bound states like the $\phi + n + n$ and/or the $\phi + p + p$. The results in this direction are presented below.

2. The role of strange sea-quarks in nucleons.

   The main quark configuration of a ϕ-meson is $s\bar{s}$. On the other hand, there are many indications on the influence of the strange sea-quark component on the nucleon wave function. It means that due to exchange of strange quarks the ϕ-meson can serve as a tool in the study of the properties of the strange component of the nuclear wave function.

3. There is a large number of experiments devoted to photoproduction and hadroproduction of ϕ-mesons where our results can be used to interpret experimental data.

In order to calculate binding energies of a few-body system, one should solve the Faddeev equations in the differential form.
The following input was used in our calculations.
The $\phi - N$ interaction was taken in the form of the Yukawa potential with the depth $\alpha = 1.25$ supporting the binding in the ϕN system with the binding energy equal to 9 MeV.
For the np triplet s-wave interaction the Malfliet-Tjon (MT) potential was used. The nn singlet s-wave interaction is based on the MT potential with a slight modification of parameters reproducing the experimental value of the nn-scattering length.
The energies of the system were obtained by solving the system of integro-differential equations by the discretization of variables [1]. The binding energy of the system ϕnn with value $E_{\phi nn} = -21.8$ MeV was obtained and the value $E_{\phi np} = -37.9$ MeV for the binding of the ϕnp system with an np pair in a triplet state.
It should be noted that the calculated binding energy of the ϕnp system is large enough to close two main ϕ-meson decay channels into K-mesons.
The dependence of the ϕnn binding energy on the parameter α of the $\phi - N$ interaction was investigated. Figure 1 shows that excited states appear in the system.
As it is seen from the results, the binding in the ϕNN system is possible even at weaker $\phi - N$ attraction in comparison with the potential used in our calculation.
It is also interesting to look whether the systems with a number of neutrons larger than two, e.g. the four-body system $\phi + 3n$, exist or not. To this aim, the folding model was used. The folding potential including the p-wave centrifugal barrier is shown in Fig. 2. This potential does not provide any binding in the system. However, having in mind that

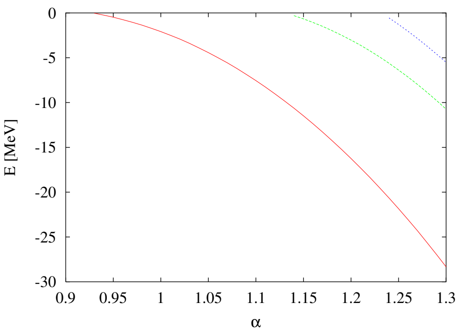

Figure 1: The dependence of the binding energy of the ϕnn system on the parameter α of the $\phi - N$ interaction.

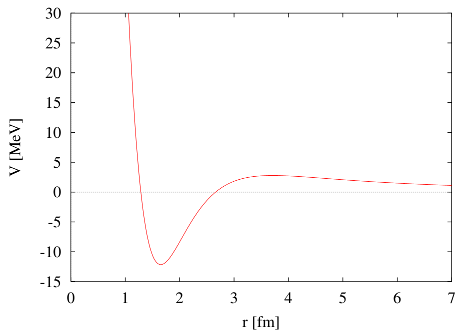

Figure 2: The folding potential with the p-wave centrifugal barrier for the four-body system $(\phi nn) + n$.

the folding model usually underestimates the binding energy of the system, the question about existence of the $\phi + 3n$ system still remains open.
Below we present the result of calculations for the systems with two ϕ-mesons like $\phi\phi N$ obtained in the framework of the Faddeev differential equations with $V_{\phi\phi}$ acting in the d-wave state [2].
The parameters of the potential $V_{\phi\phi}$ are chosen to fit (together with the centrifugal barrier) the position and width of the $f_2(2010)$- resonance which has one mode of decay into two ϕ-mesons.

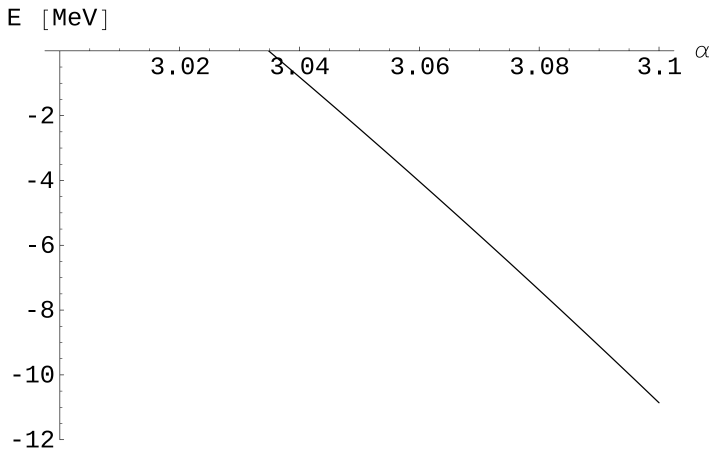

Figure 3: The dependence of the binding energy of the $\phi\phi n$ system on the parameter α of the $\phi - N$ interaction.

Some approximations were made in the calculations. For example, it was found reasonable to suppose the line configuration of the $\phi\phi N$ system as the most probable one, due to the strong repulsion of two ϕ mesons, being in the D-state.
The dependence of the energy of the system $\phi\phi N$ on the depth of the ϕN potential α is shown in Fig. 3. One can see that the system becomes the bound one only at $\alpha = 3.035$ which is quite large in comparison with the initial value.
The study of the ϕ-meson-nuclear systems may also shed light on the distributions of s and $\bar{s}$ sea quarks in nuclei (as well as in nucleon immersed in nuclear medium) and on the

possible appearance of many-body effects related to the exchange of sea s and $\bar{s}$ quarks belonging to different baryons.

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CALCULATIONS OF THE K^{+}-NUCLEUS MICROSCOPIC OPTICAL POTENTIAL AND THE CORRESPONDING DIFFERENTIAL ELASTIC CROSS SECTIONS

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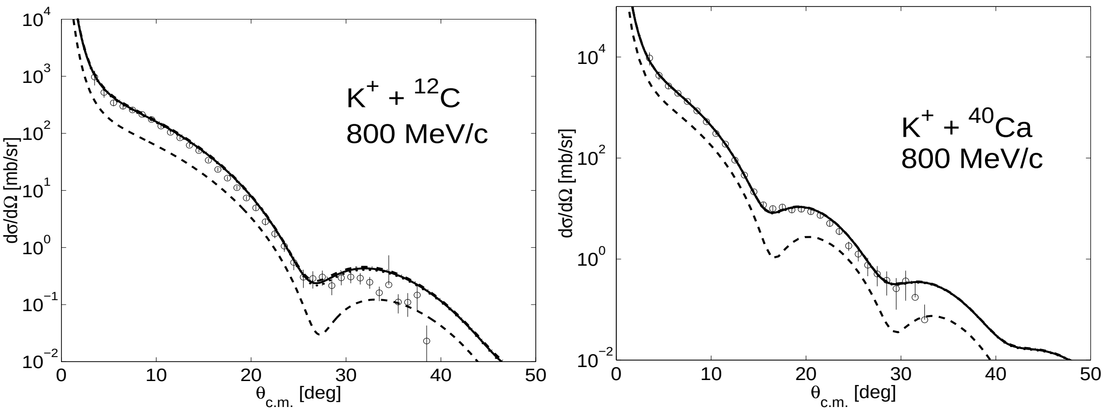
In Ref. [1], the elastic scattering cross sections of kaons K^+ with momenta 0.635, 0.715, and 0.8 GeV/c off the ^{12}C and ^{40}Ca nuclei were calculated. The microscopic optical potential (OP) derived in the high-energy approximation [2] was utilized in the form

$$U^H = V^H + iW^H = -\frac{\hbar c\,\beta}{(2\pi)^2} \sum_{\nu=p,n} \bar{\sigma}_K^\nu(\bar{\alpha}_K^\nu + i) \int_0^\infty dq\; q^2 j_0(qr)\rho_\nu(q) f_K^\nu(q). \tag{1}$$

Here the existing data were used on the KN-scattering (form factor f_K, total cross section $\bar{\sigma}$, ratio $\bar{\alpha}_K = ReF_{KN}(0)/ImF_{KN}(0)$) and nuclear densities $\rho(r)$. In view of the relations $E \gg m^K$, U^H the ordinary Klein-Gordon-Fock equation can be transformed to

$$(\Delta + k^2)\psi(\mathbf{r}) \;=\; 2\mu\gamma^{(r)}\Big(U - \frac{U^2}{2E}\Big)\psi(\mathbf{r}), \qquad U \;=\; U^H + U_C \tag{2}$$

with a relativistic momentum k. The factor $\gamma^{(r)}$ is the ratio of reduced relativistic energy to the non-relativistic reduced mass μ and is presented in different works in slightly different forms caused by some additional approximations. The Figure exhibits differential



cross sections calculated by using Eqs. (1)-(2). The dashed curves correspond to the case without the relativistic transform of OP ($\gamma^{(r)}$=1), while the solid curves demonstrate a significant effect of relativization when the respective factors calculated for all afore-said cases turned out to be in the limits of $\gamma^{(r)} \simeq 1.56 \doteq 1.87$. At the same time, it was demonstrated in [1] that the different methods of relativization result in almost the same cross sections. Also, a small difference between the densities from a number of tables do not provide a noticeable effect. Besides, the effect of the $U^2/2E$ term on the cross sections is rather week, too. As to the total reaction cross sections, the relativization increases them by about 30%, and some additional enhancement takes place when a surface absorption term ($\simeq dW/dr$) is included in OP.

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RELATIVISTIC DESCRIPTION OF THE DEUTERON WITHIN THE BETHE-SALPETER APPROACH

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The deuteron is an object of intensive investigations as the simplest bound neutron-proton system. Throughout more than 40 years many methods for the description of the deuteron have been elaborated. Using a separable form of interaction allows one to simplify calculations. That is why there are separable approximations intended only to reproduce the behavior of the corresponding realistic potentials and to be used in calculations instead of more complicated originals (see for details, e.g., [1]).
This idea was developed in [2-5] to describe uncoupled partial-wave states in the elastic np scattering for $T_{\rm Lab}$ up to 3 GeV.
The rank-six separable interaction kernel for the triplet partial-wave state ${}^3S_1^+$-${}^3D_1^+$ was proposed in our work [6]. It is a continuation of the previous one [5] where the uncoupled partial-wave states with the total angular momenta $J = 0, 1$ were considered. Various deuteron characteristics were investigated using the elaborated kernel. Parameters of the model were obtained from the fitting of experimental data for the phase shifts (the SAID program: http://gwdac.phys.gwu.edu) and low-energy characteristics.
The calculated low-energy scattering parameters and deuteron characteristics compared with the corresponding experimental values and other models can be found in [6].

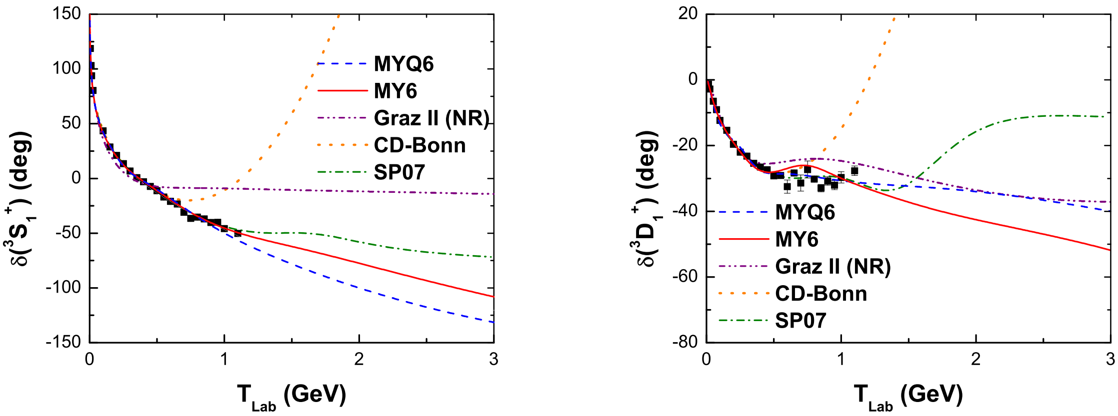

Figure 1: The model phase shifts for the ${}^3S_1^+$ and ${}^3D_1^+$ partial-wave states are shown. Two relativistic separable kernel cases MY6 and MYQ6 are compared to those of Graz II (NR) [8], CD-Bonn [11] and of the empirical SP07 SAID solution [7].

In Fig. 1, the obtained phase shift for the ${}^3S_1^+$ and ${}^3D_1^+$ partial-wave states are compared with experimental data and, in addition to the afore-said theoretical models, with the empirical SP07 SAID solution [7]. The nonrelativistic Graz II [8] potential (Graz II (NR)) is considered here as an alternative separable model. The relativistic interaction kernel

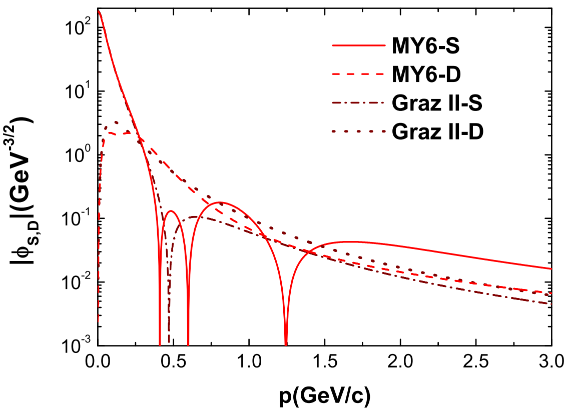

Figure 2: Wave functions $\phi(\bar{p}_0, \mathbf{p})$ [5] for the ${}^3S_1^+$ and ${}^3D_1^+$ partial-wave states at $p_0 = M_d/2 - E_{\mathbf{p}}$ are presented. They are written in the deuteron rest frame. The MY6 model [6] (MY6-S red solid line corresponds to ${}^3S_1^+$ partial-wave state, MY6-D red dashed line - to ${}^3D_1^+$) is compared with Graz II [9] (Graz II-S brown dash-dotted line - ${}^3S_1^+$ wave function, Graz II-D brown dotted line - ${}^3D_1^+$ wave function).

Graz II [9] is not presented because in this case the calculation of the phase shifts cannot be performed in the whole energy range where experimental data are available. As it was discussed in [1, 5], if the Graz II model is used, it is impossible to perform calculations in principle when T_{Lab} exceeds some limit value (which depends on the parameters of separable model functions), whereas our aim is to compare our MY6 and MYQ6 results with other models for all available experimental data. As one can see from Fig. 1, the ${}^3S_1^+$ phase shifts are well described by both MY6 and MYQ6 parameterizations. The Graz II (NR) works at $T_{\mathrm{Lab}} \leqslant 0.4$ GeV. For the ${}^3D_1^+$ state MY6 and MYQ6 also provide a good description. The Graz II (NR) shows only some correspondence with the data at $T_{\mathrm{Lab}} \leqslant 0.4$ GeV. SP07 is good for all experimental data. CD-Bonn was constructed for $T_{\mathrm{Lab}} \leqslant 350$ MeV and is perfect within this T_{Lab} interval. However, its high-energy behavior means that other models should be used at $T_{\mathrm{Lab}} > 500$ MeV, whereas the interaction kernels MY6 and MYQ6 demonstrate a reasonable behavior in the whole energy range.
As any other phenomenological model, ours can describe on-shell characteristics quite easily. However, when the coupled ${}^3S_1^+$-${}^3D_1^+$ channel is considered, phase shifts and low-energy characteristics are not the only observables which must be described. It is also important to look at properties of the deuteron BS amplitude (wave function). Therefore, in calculations we take into account that the obtained ${}^3S_1^+$ and ${}^3D_1^+$ wave functions $\phi(\bar{p}_0, \mathbf{p})$ at $\bar{p}_0 = M_d/2 - E_{\mathbf{p}}$ [10] (Fig. 2) should be similar to other discussed models in the energy region where their properties play a key role. The relativistic Graz II model is presented for comparison.
The proposed separable models MY6 and MYQ6 can be used in calculations of various reactions with the deuteron, e.g., the deuteron photo- and electrodisintegration etc. Additional parameters α provide integrands containing separable model functions to have no poles at any $\mathbf{p}$. Therefore, functions of this type allow one to perform numerical calculations of the electrodisintegration far from the threshold without resorting quasipo-

tential or nonrelativistic approximations. Comparing with other separable and realistic potential models, we can demonstrate the merits of separable kernels with the α-modified model functions. The CD-Bonn potential, which was constructed for $T_{\rm Lab} \leqslant 350\,\mathrm{MeV}$ and works in this energy interval very well, cannot just be simply extrapolated to higher energies. The Graz II interaction kernel is useless in high-energy calculations because they are impossible in principle in that case. On the contrary, our model has no these limitations. New experimental data for exclusive electrodisintegration of the deuteron at high momentum transfer [12] can be a good instrument for testing the proposed relativistic models. The specific arrangement of the experiment, when the final state interaction (FSI) effects are minimized, allows one to compare the results obtained within the plane wave impulse approximation (PWIA). Therefore, it is a chance to investigate the influence of nucleon momentum distributions produced by various models of NN interactions.
In our work [13], the exclusive cross section of electrodisintegration $d^2\sigma/(dQ^2 d|\mathbf{p}_n|)$ for kinematic conditions of the Jefferson Laboratory experiment [12] was calculated within the Bethe-Salpeter approach with the rank-six separable kernel MY6 [6]. The calculations were performed within the relativistic PWIA. The obtained results were compared with the experimental data and two theoretical models, the nonrelativistic Graz II (NR) [8] and relativistic Graz II [9] separable interaction kernels.

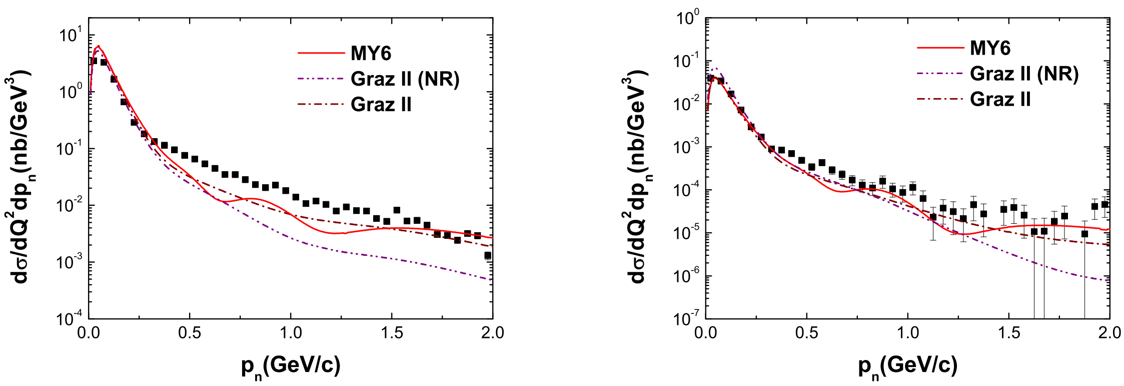

Figure 3: The cross section $d^2\sigma/(dQ^2 d|\mathbf{p}_n|)$ [13] depending on neutron momentum $|\mathbf{p}_n|$ is considered for $Q^2 = 2 \pm 0.25\,\mathrm{GeV}^2$ (left panel) and $Q^2 = 5 \pm 0.5\,\mathrm{GeV}^2$ (right panel). Calculations with the Graz II (NR) [8] (purple dash-dot-dotted line), Graz II [9] (brown dash-dotted line) and MY6 [6] (red solid line) models are present. The dipole fit model [14] for nucleon form factors is used.

Figure 3 illustrates the cross section depending on outgoing neutron momentum $|\mathbf{p}_n|$ for transfer momenta Q^2=2 GeV^2 and Q^2=5 GeV^2. The dipole fit model [14] for the nucleon electromagnetic form factors was used. One nonrelativistic Graz II (NR) and two relativistic MY6, Graz II separable kernels of NN interactions were investigated. Good agreement with the experimental data at low neutron momenta $|\mathbf{p}_n| < 0.25\,\mathrm{GeV/c}$ can be seen in the figures. The discrepancy between the theoretical models and the experimental data increases with $|\mathbf{p}_n| > 0.25\,\mathrm{GeV/c}$ for all the considered models. However, we see the agreement of the relativistic models (MY6, Graz II) with the experimental data at high neutron momenta. Moreover, the relativistic description becomes better with Q^2 increasing, and theoretical curves go practically along experimental points at $Q^2 = 5\,\mathrm{GeV}^2$. Therefore, relativistic effects play an important role in the description of the deuteron electrodisintegration at high momentum transfer and high neutron momenta.

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VISCOSITY OF HADRON MATTER WITHIN A RELATIVISTIC MEAN-FIELD MODEL WITH SCALED HADRON MASSES AND COUPLINGS

V. D. Toneev and A. S. Khvorostukhin

In the past, transport coefficients for the nuclear matter were actively studied (see Introduction in [1]). Recently, interest in the transport coefficient issue sharply increased in heavy-ion collision physics. Values of the elliptic flow v_2 observed at RHIC proved to be larger than v_2 at SPS. This finding is interpreted as that a quark-gluon plasma (QGP) created at RHIC behaves as a nearly perfect fluid with a small value of the shear viscosity-to-entropy density ratio, η/s. The latter statement was confirmed by non-ideal hydrodynamic analysis of these data. Thereby, it was claimed that a new state produced at high temperatures is most likely not a weakly interacting QGP, as it was originally assumed but a strongly interacting QGP. The interest was also supported by a new theoretical perspective, namely, $\mathcal{N} = 4$ supersymmetric Yang-Mills gauge theory using the Anti de-Sitter space/Conformal Field Theory (AdS/CFT) duality conjecture. Calculations in this strongly coupled theory demonstrate that there is a minimum in the η/s ratio: $\eta/s \approx 1/(4\pi)$. It was thereby conjectured that this relation is in fact a lower bound for the specific shear viscosity in all systems and that the minimum is reached in the hadron-quark transition critical point (at $T = T_c$).

Constructing the Equation of State (EoS) we assume here a relevance of the (partial) chiral symmetry restoration at high baryon densities and/or temperatures manifesting themselves in the form of the Brown-Rho scaling hypothesis: masses and coupling constants of all hadrons decrease with a density increase in approximately the same way. Simultaneously, our model fulfills various constraints known from analysis of atomic nuclei, neutron stars and HIC.

Within our relativistic mean-field model with Scaled Hadron Masses and Couplings (SHMC) [3] we present the Lagrangian density of the hadronic matter as a sum of several terms:

$$\mathcal{L} = \mathcal{L}_{\rm bar} + \mathcal{L}_{\rm MF} + \mathcal{L}_{\rm ex} \ . \tag{1}$$

The Lagrangian density of the baryon component interacting via σ, ω mean fields is as follows:

$$\mathcal{L}_{\rm bar} \ = \sum_{b\in\{\rm bar\}} \left[i\bar{\Psi}_b \left(\partial_\mu + i\, g_{\omega b}\, \chi_\omega\ \omega_\mu\right)\gamma^\mu\, \Psi_b - m_b^*\, \bar{\Psi}_b\, \Psi_b \right]. \tag{2}$$

The considered baryon set is $\{b\} = N(938)$, $\Delta(1232)$, $\Lambda(1116)$, $\Sigma(1193)$, $\Xi(1318)$, $\Sigma^*(1385)$, $\Xi^*(1530)$, and $\Omega(1672)$, including antiparticles. The used σ-field dependent effective masses of baryons are

$$m_b^*/m_b = \Phi_b(\chi_\sigma \sigma) = 1 - g_{\sigma b}\ \chi_\sigma\ \sigma/m_b\,,\ b \in \{b\}\ . \tag{3}$$

In Eqs. (2), (3) $g_{\sigma b}$ and $g_{\omega b}$ are the coupling constants and $\chi_\sigma(\sigma)$, $\chi_\omega(\sigma)$ are the coupling scaling functions.

The σ- and ω-meson mean field contribution is given by

$$\mathcal{L}_{\rm MF} = \frac{\partial^\mu \sigma \; \partial_\mu \sigma}{2} - \frac{m_\sigma^{*2}\, \sigma^2}{2} - U(\chi_\sigma \sigma) - \frac{\omega_{\mu\nu}\, \omega^{\mu\nu}}{4} + \frac{m_\omega^{*2}\, \omega_\mu \omega^\mu}{2} \;, \tag{4}$$
$$\omega_{\mu\nu} = \partial_\mu \omega_\nu - \partial_\nu \omega_\mu \;, \quad U(\chi_\sigma \sigma) = m_N^4 (\frac{b}{3} f^3 + \frac{c}{4} f^4), \quad f = g_{\sigma N}\; \chi_\sigma \; \sigma / m_N \,.$$

There exist only σ and ω_0 mean field solutions of the equations of motion. The mass terms of the mean fields are

$$m_m^* / m_m \;\; = \;\; |\Phi_m(\chi_\sigma \sigma)| \,, \quad \{m\} = \sigma, \omega \,. \tag{5}$$

The dimensionless scaling functions Φ_b and Φ_m, as well as the coupling scaling functions χ_m, depend on the scalar field in a combination $\chi_\sigma(\sigma)\;\sigma$. We assume the approximate validity of the Brown-Rho scaling ansatz in the simplest form

$$\Phi = \Phi_N = \Phi_\sigma = \Phi_\omega = \Phi_\rho = 1 - f. \tag{6}$$

The third term in the Lagrangian density (1) includes meson quasiparticle excitations: $\pi; K, \bar{K}; \eta(547); \sigma', \omega', \rho'; K^{*\pm,0}(892), \eta'(958), \phi(1020)$. The parameters of the relativistic mean field model, C_σ, C_ω and the self-interaction potential U are adjusted to reproduce the nuclear matter properties at the saturation for $T = 0$. The choice of parameters and other details of the SHMC model can be found in [2, 3]. One can demonstrate [2, 3] that the SHMC model describes the nucleon optical potential $U_{\rm opt}$ in an optimal way and the pressure at $T = 0$ calculated in the SHMC model well satisfies the experimental constraints coming from the analysis of an elliptic flow.
Within the SHMC model we calculate different thermodynamic quantities in thermal equilibrium hadron matter at fixed temperature T and baryon chemical potential $\mu_{\rm bar}$. To extend this hadronic EoS to higher temperatures, we use the Heavy Quark Bag (HQB) model. This two-phase model includes the first order phase transition and is in agreement with the lattice data for the pressure and energy density [2, 3].
The developed EoS was applied to study the kinetic coefficients. In the relaxation time approximation we derived expressions for the shear (η) and bulk (ζ) viscosities in the case when the quasiparticle depends on the temperature via the dispersion relation $E(\vec{p}) = \sqrt{\vec{p}^{\,2} + m^{*\,2}(T, \mu)}$ and on the mean fields.
Finally, we obtain expressions for the shear and bulk viscosity as follows [1]:

$$\eta = \frac{1}{15T} \sum_a \int d\Gamma \, \tau_a \frac{\vec{p}_a^{\,4}}{E_a^2} \; [F_a^{\rm eq} \; (1 \mp F_a^{\rm eq})] \,, \tag{7}$$

$$\zeta = -\frac{1}{3T} \sum_a \int {\rm d}\Gamma \, \tau_a \; \frac{\vec{p}_a^{\,2}}{E_a} \; F_a^{\rm eq} \; (1 \mp F_a^{\rm eq}) \, Q_a \tag{8}$$

with the factor

$$Q_a = -\Bigg\{ \frac{\vec{p}_a^{\,2}}{3E_a} + \left(\frac{\partial P}{\partial n_{\rm bar}} \right)_{\epsilon, n_{\rm str}} \left[\frac{\partial (E_a + X_a^0)}{\partial \mu_{\rm bar}} - t_b^{\rm bar} \right] + \left(\frac{\partial P}{\partial n_{\rm str}} \right)_{\epsilon, n_{\rm bar}} \left[\frac{\partial (E_a + X_a^0)}{\partial \mu_{\rm str}} - t_a^{\rm str} \right]$$
$$- \left(\frac{\partial P}{\partial \epsilon} \right)_{n_{\rm bar}, n_{\rm str}} \times \left[E_a + X_a^0 - T \frac{\partial (E_a + X_a^0)}{\partial T} - \mu_{\rm bar} \frac{\partial (E_a + X_a^0)}{\partial \mu_{\rm bar}} - \mu_{\rm str} \frac{\partial (E_a + X_a^0)}{\partial \mu_{\rm str}} \right] \Bigg\} \,,$$

the equilibrium distribution function F^{eq}_a and the relaxation time of the species a

$$\tau_a^{-1}(T,\mu) = \sum_{a'} n_{a'}(T,\mu) \left\langle v_{aa'} \sigma^t_{aa'}(v_{aa'}) \right\rangle . \tag{9}$$

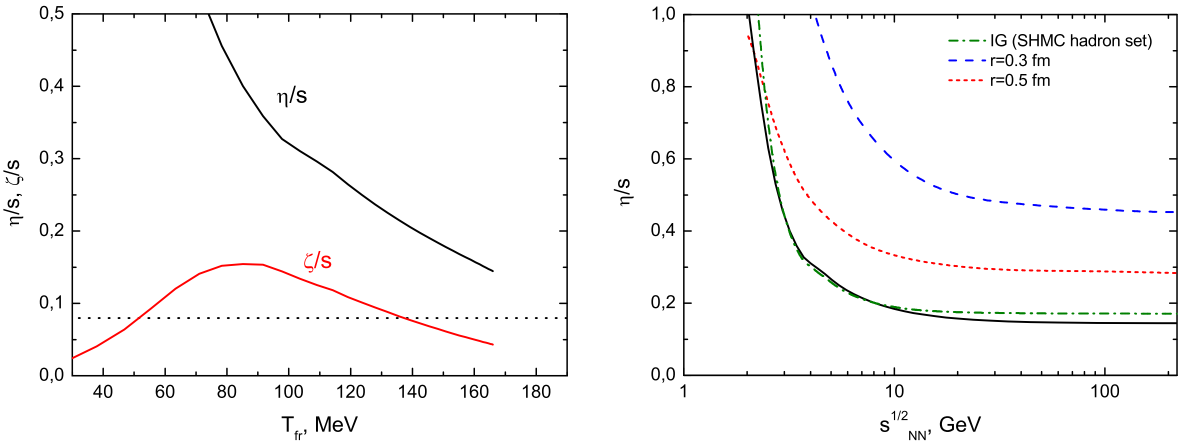

Figure 1: Predictions of the SHMC model (solid lines) for specific shear and bulk viscosities calculated for central Au+Au collisions along the freeze-out curve as a function of the freeze-out temperature (left panel) and the collision energy $s^{1/2}_{NN}$ (right) for the baryon enriched system. The dotted line is the lower AdS/CFT bound $\eta/s = 1/4\pi$. The dashed and short-dashed curves are the results of the excluded volume hadron gas model. The dot-dashed line corresponds to the IG model with the same set of hadrons as for the SHMC model.

In Fig. 1, the model predictions for the reduced share and bulk viscosity are presented at the freeze-out. The η/s ratio decreases monotonously with increase of the temperature, being higher than the lower bound $1/4\pi$ but tending to it with further increase in the freeze-out temperature $T_{\rm fr}$. The value ζ/s exhibits a maximum at $T_{\rm fr} \sim 85$ MeV and then goes to zero with a subsequent increase of $T_{\rm fr}$. At $T \gtrsim 100$ MeV the values of shear and bulk viscosities become rather close $(\eta/s)_{\rm fr} \simeq 2(\zeta/s)_{\rm fr}$.
The case with the phase transition is illustrated in Fig. 2 and compared with the results of the NJL model. Below the critical temperature T_c the ratios are the same, as shown above in Fig. 1. For $T > T_c$ the original NJL model gives a continuous smooth line (a crossover) for $\mu_{\rm bar} = 0$. In our two-phase SHMC-HQB model there is a jump at T_c in both the η/s and ζ/s ratios. This jump is a particular property of the first order phase transition.
Concluding, the modified relativistic mean-field σ-ω-ρ model with scaled hadron masses and couplings was generalized to finite temperatures. The EoS for $T = 0$ satisfies general constraints known from atomic nuclei, neutron stars and those coming from the flow analysis of HIC data. The developed SHMC model in combination with the heavy quark bag model, which describes the EoS of hot and dense hadronic matter in a broad range of temperatures and baryon densities, is applied for describing kinetic coefficients. The predictions for (η/s) and (ζ/s) at the freeze-out surface are made as well as particularities of the share and bulk viscosities near the phase transition point are demonstrated.

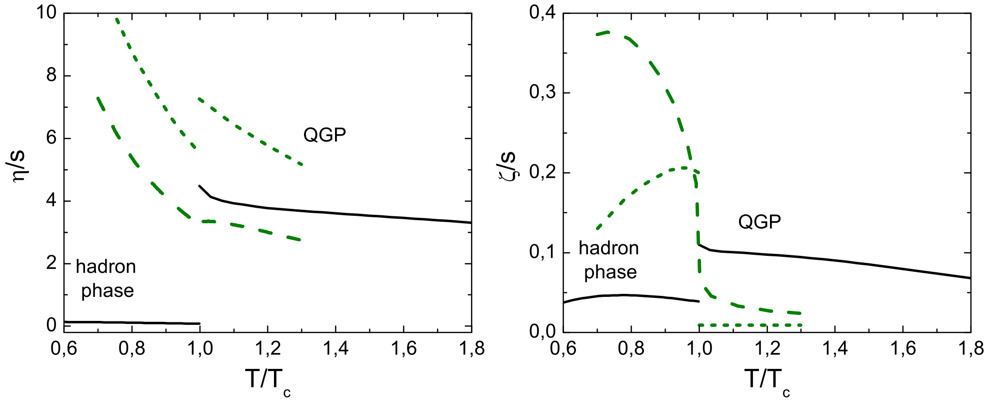

Figure 2: The T-dependence of the shear (left panel) and bulk (right panel) specific viscosities within our two-phase SHMC-HQB model (solid lines) for $\mu_{\rm bar}^{\rm CEP}$ =990 MeV, corresponding to the critical end point. The NJL model results for $\mu_{\rm bar}^{\rm CEP}$ and for $\mu_{\rm bar}$ slightly above $\mu_{\rm bar}^{\rm CEP}$ are plotted by the long-dash and short-dash lines, respectively.

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CALCULATION OF THE CROSS SECTION AND THE TRANSVERSE-LONGITUDINAL ASYMMETRY OF THE PROCESS ^{3}He$(e, e'p)pn$ AT MEDIUM ENERGIES WITHIN THE UNFACTORIZED GENERALIZED GLAUBER APPROACH

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The cross section and the transverse-longitudinal asymmetry A_{TL} of the three-body-break-up process $^3He(e, e'p)pn$ have been calculated by an unfactorized and parameter-free approach based upon realistic few-body wave functions corresponding to the $AV18$ interaction, treating the rescattering of the struck nucleon within a generalized eikonal approximation. The results of calculations exhibit good agreement with the recent JLab experimental data and show the dominant role played by the Final State Interaction which, however, in the region of missing momentum, $300 \lesssim p_m \lesssim 600\ MeV/c$, and removal energy corresponding to the two-body kinematics peak and higher, $E_m \gtrsim p_m^2/4m_N$, is dominated by single nucleon rescattering, providing evidence that the final state interaction is mainly due to the one between the struck nucleon and a nearby correlated one. The cross section of the process has the form

$$\frac{d^6\sigma}{d\Omega' dE'\ d\Omega_{p_1} dE_{p_1}} = \left| \frac{\mathbf{p}_1^2}{\dfrac{|p_1|}{E_{p_1}} + \dfrac{|\mathbf{p}_1| - |\mathbf{q}|\cos\theta}{E^*_{A-1}}} \right| \frac{dE_m}{dE_{p_1}}\ \sigma_{Mott} \sum_i\ V_i\ W_i^A(\nu, Q^2, \underline{\mathrm{p}}_m, E_m), \quad (1)$$

where $i \equiv \{L, T, TL, TT\}$, V_i are kinematical factors, and the nuclear structure functions W_i^A result from proper combinations of the polarization vector of the virtual photon ε_λ^μ and the hadronic tensor $W_{\mu\nu}^A$ the latter depending upon the nuclear current operators $\hat{J}_\mu^A(0)$ which, besides the electromagnetic part of the interaction, describes also the process of multiple elastic scattering in the final state. In the present approach, the rescattering process is treated within the Generalized Eikonal Approximation, according to which the corresponding matrix elements $\hat{J}_\mu^A(0)$ read as

$$\begin{aligned} J_\mu^3 &= \sum_\lambda \int \frac{d\mathbf{p}}{(2\pi)^3}\frac{d\kappa}{(2\pi)^3} S_\Delta^{FSI}(\mathbf{p},\kappa)\langle\lambda_\mathbf{f}|\mathbf{j}_\mu(\kappa - \mathbf{p_m};\mathbf{q})|\lambda\rangle\mathcal{O}(\mathbf{p_m} - \kappa, \mathbf{p}, \mathbf{k_{rel}}; \mathcal{M}_\mathbf{3}, \mathbf{S_f}, \sigma_\mathbf{f}, \lambda) \\ &= J_\mu^{3(PWIA)} + J_\mu^{3(1)} + J_\mu^{3(2)}, \end{aligned} \quad (2)$$

where S_Δ^{FSI} is the rescattering S-matrix within the Eikonal Approximation, $\langle\lambda_f|j_\mu(\kappa - \mathbf{p}_m; \mathbf{q})|\lambda\rangle$ is the nucleonic electromagnetic current and $\mathcal{O}$ is the nuclear overlap in momentum space

$$\mathcal{O}(\mathbf{p}_m - \kappa, \mathbf{p}, \mathbf{k_{rel}}\,; \mathcal{M}_\mathbf{3}, \mathbf{S_f}, \sigma_\mathbf{f}, \lambda) = \int \mathbf{d}\rho\mathbf{dr}\mathbf{e}^{\mathbf{i}(\mathbf{p_m}-\kappa)\,\rho}\mathbf{e}^{\mathbf{ip\,r}}\mathbf{\Phi}_{\frac{1}{2},\mathcal{M}_\mathbf{3}}(\rho,\mathbf{r})\mathbf{\Phi}^{\mathbf{k_{rel}}*}_{\mathbf{S_f},\sigma_\mathbf{f}}(\mathbf{r})\chi^\dagger_{\frac{1}{2}\lambda} \quad (3)$$

The left-right asymmetry is defined by

$$A_{TL} = \frac{d\sigma(\phi = 0^o) - d\sigma(\phi = 180^o)}{d\sigma(\phi = 0^o) + d\sigma(\phi = 180^o)}. \quad (4)$$

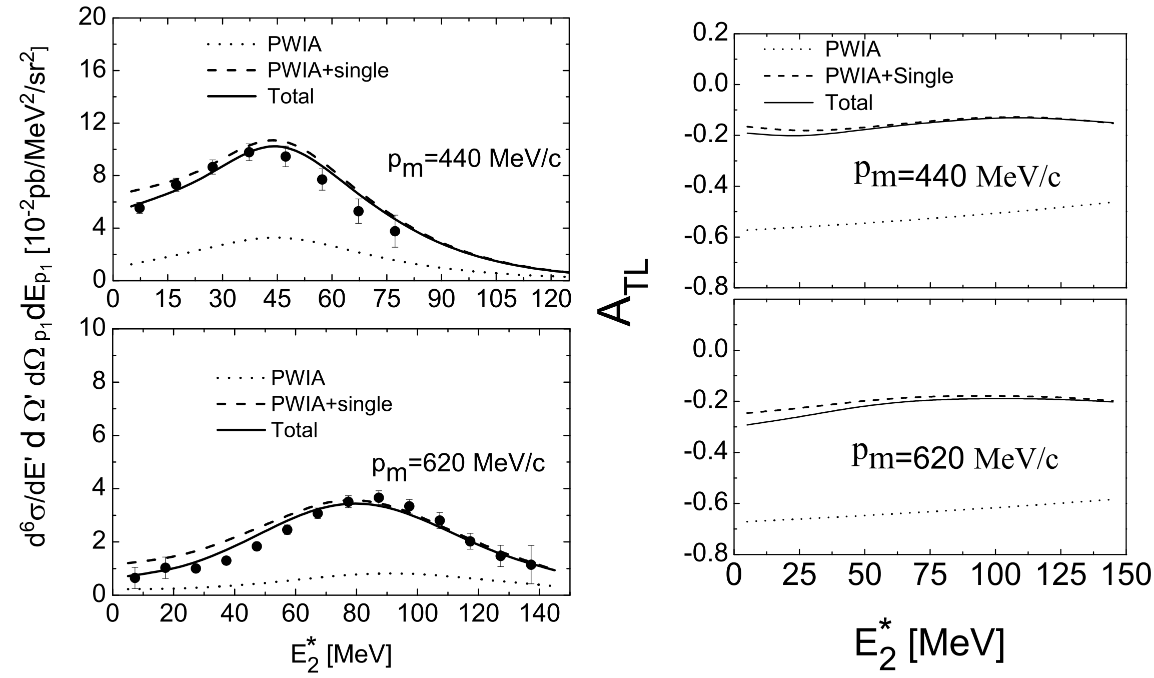

Figure 1: The differential cross section Eq. (1) (left panels) and the left-right asymmetry A_{TL} Eq. (4) (right panels) for the process $^3\mathrm{He}(e, e'p)pn$ calculated at two values of the missing momentum *vs.* the excitation energy of the two-body final state $E_2^* = E_m - E_{min}$ ($E_{min} = 2m_p + m_n - M_3$). Dotted lines: PWIA approximation; dashed and solid lines: unfactorized calculations with single and double rescattering in the final state, respectively. Experimental data from [2].

The results of our calculations of the cross section (1) and predictions for the left-right asymmetry A_{TL} (4) are presented in Fig. 1 (details of calculations and discussions can be found in Ref. [1]).

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DIMUON PRODUCTION BY LASER-WAKEFIELD ACCELERATED ELECTRONS

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The possibility to produce strong electric fields of the order of 10−100 GV/m with present laser facilities is a great advantage for laser-wakefield accelerators [1], which allows, in principle, to construct compact accelerating devices for particle and nuclear physics. The successful production of high-quality electron beams in such laser-driven accelerators with electron energies of the order of 1 GeV was reported recently [2, 3]. Electron beams with energies exceeding 1 GeV are interesting for many applications in particle and nuclear physics [4]. One of the attractive subjects is related to the neutrino physics. For example, for studying neutrino oscillations it would be important to have two types of neutrinos with fixed intensity. This may be obtained in muon decays $\mu^+ \to e^+ + \nu_e + \bar{\nu}_\mu$ and $\mu^- \to e^- + \bar{\nu}_e + \nu_\mu$, where muon and electron neutrinos (or antineutrinos) are produced in equal parts. Therefore, it is interesting to estimate whether the high-energy laser-driven electrons can produce a sizeable amount of muon pairs for future applications. Together with neutrino oscillation, such high-intensity muon sources may be used in studying other fundamental problems of lepton physics, say the search for lepton flavor violation and the measurement of the muon anomalous magnetic moment. The aim of the present study is to analyze the possibility of muon pair creation in the interaction of high-energy laser driven electrons within a heavy (high-Z) target in a table top configuration [5].

The electromagnetic sources of the $\mu^+\mu^-$ (dimuon) production are described by the following elementary processes:

$$\gamma + A \to A + \mu^+\mu^- \tag{1}$$

and

$$e + A \to e' + A + \mu^+\mu^- \tag{2}$$

In the first case, the dimuons are produced in the interaction of real (bremsstrahlung) photons within the electric field of the high-Z target nuclei. In the second case, the dimuons are produced in the interaction of high-energy electrons with nuclei. First, we analyzed different aspects of dimuon productions in elementary processes and then we used them for evaluation of their total yield for given electron beam and target properties. For the former ones we use the conditions of electron beams, as reported for the laser-wakefield accelerator in Ref. [2]. The electron energy is about 0.5−1 GeV and the electron flux is about 20 pC which corresponds to $N_0^e \simeq 1.248 \times 10^8$ electrons in a bunch. In our estimates we assume the same flux for electron energies up to 10 GeV. We consider a gold target with thickness of L=0.1−1 cm. We evaluate the di-muon yield using transport-kinetic like model [5].

In Fig. 1, we present the total dimuon yield in interactions of relativistic electrons with a gold target which is a sum of the (1) and (2) contributions as a function of the primary electron energy E_e^0 and target thickness L.

This result illustrates the effectiveness of the laser wakefield accelerator, impinging on a thick high-Z target as a source of muon pairs. For a 1 cm thick gold target, 1.25×10^8 electrons in a 20 pC bunch with energy of 1 (10) GeV in the initial state produce about 1×10^2 (5×10^3) dimuons with pair energies centered at 1 GeV. To get 10^6 dimuons from

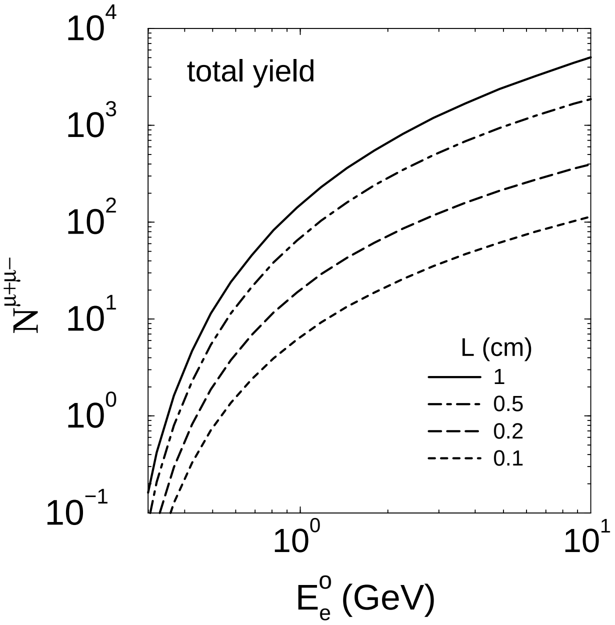

Figure 1: The total yield of dimuons in interactions of relativistic electrons with the gold target as a function of the primary electron energy E_e^0 and target thickness L.

the muon factory, one needs $10^{10} - 10^{11}$ primary electrons in a bunch. Such intensities with power of 100 J seem to be quite realistic in near future, requiring ultra-high intensity laser pulses with efficient acceleration mechanisms. Thus, the configuration of a laser driven electron accelerator and thick high-Z target may serve as an all-optics table top device for muon pair production. Their advantage at present is related to the high density of the particles, the excellent normalized emittance, the small size of the driver, possible high power scalability, synergies with nuclear fusion, etc. The produced muons may be used in studying various aspects of muon and neutrino physics and to be considered as an important step towards investigations of more complicated electron induced elementary processes.

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3. N. V. Antonenko (principal investigator), R. V. Jolos, G. G. Adamian, A. K. Nasirov, Sh. A. Kalandarov, V. V. Sargsyan, A. S. Zubov, "Dynamics of formation of heaviest nuclei and their structure", RFBR-China, 08-02-92202-GFEN-a, 2009-2010.

4. G. G. Adamian (principal investigator), N. V. Antonenko, A. K. Nasirov, A. V. Andreev, Sh. A. Kalandarov, S. N. Kuklin, V. V. Sargsyan, T. M. Shneidman, A. S. Zubov, "Non-equilibrium and equilibrium phenomena in production and structure of superheavy and exotic nuclei", 10-02-91340-NNIO-a, 2010-2011.

5. S. N. Ershov, "Theoretical investigation of the drip-line nuclei: structure and nuclear reactions" RFBR 08-02-00892, 2008-2010; N. Shulgina (Russian Research Center "Kurchatov Institute", Principal Investigator) 4 participants.

6. "Spectral and evolutional problems", RFBR 09-01-90408, 2009-2010 (Principal Investigator: A. A. Shkalikov, Moscow State University; participant from JINR: A. K. Motovilov).

7. "Operator models in mathematical physics and their spectral analysis", RFBR 10-01-00423, 2010-2012 (Principal Investigator: A. A. Shkalikov, Moscow State University; participant from JINR: A. K. Motovilov).

8. "Dynamics of coherent systems in optical lattices and traps", RFBR 08-02-00118 (Russia), 12.2008 - 12.2010 V.I. Yukalov (Principal Investigator), 6 participants (R.G. Nazmitdinov, V. O. Nesterenko,).

9. A. N. Sissakian, A. S. Sorin, V. D. Toneev, ..., "Search for a mixed phase of strongly interacting nuclear matter", RFBR 08-02-010001-a, 01.2008-12.2010.

10. A. Sushkov, "Investigation of collective states in the well-deformed nucleus ^{160}Dy by the precision spectroscopy methods and their theoretical analysis", RFBR 08-02-00622, 2008-2010, V. G. Kalinnikov (DLNP, Principal Investigator).

11. A. S. Zubov, grant of the President of Russian Federation for the young PhD scientists MK-2429.2010.2; 05.2010-11.2011

DFG grants

1. “Subspace perturbation problem and operator Riccati equation”, DFG-RFBR 436 RUS 113/ 817 (06–01–04003), 2005–2010 (Principal Investigators: S. Albeverio, Institute for Applied Mathematics, Bonn University, Bonn, Germany, and A. K. Motovilov)

2. “Interaction of different mesons with the three- and four-nucleon systems”, DFG 436 RUS 113/761, 2007–2010 (Principal Investigators: W. Sandhas, Institute for Physics, University of Bonn, Bonn, Germany, and V. B. Belyaev; participant from BLTP: I. I. Shlyk).

3. R. V. Jolos, G. G. Adamian, N. V. Antonenko, A. V. Andreev, A. K. Nasirov, S. N. Kuklin, Sh. A. Kalandarov, V. V. Sargsyan, “Dynamics of cluster formation in nuclear structure, fusion and fission”, DFG–RFBR–08–02–91961, 2008–2009.

4. “Anisotropic scattering of ultracold atoms and polar molecules in waveguides”, DFG Schm 885/18-1, 2010 (Principal Investigators: P. Schmelcher, Center of Quantum Optics, University of Hamburg, Hamburg, Germany, and V. S. Melezhik).

5. V. O. Nesterenko, DFG grant RE 322/12-1, “Investigation of nuclear dynamics with Skyrme forces” (P.-G. Reinhard); grant GSI-F+E-2010-12 (P.-G. Reinhard + GSI people)

Other grants

1. V. O. Nesterenko,”Spectroscopy of atomic clusters”, grants of University of Paul Sabatier (Toulouse, France), E. Suraud (leader); No 684 (2009); No 111 (2010).

2. R.G. Nazmitdinov, Universidad de les Illes Balears (Spain), Grant FIS2008-00781/FIS, 12.2008-12.2010; Ll. Serra (leader), 10 participants.

EDUCATIONAL ACTIVITY

1. V. V. Voronov: 2 regular lecture courses “Theory of vibrations”, “Fundamentals in physics of nuclei and elementary particles”, Tver’ State University;

   lecture course “Many-body system methods in nuclear theory”, Intern. Dubna University.

2. V. A. Kuz’min: Lecture course on ”Theory of atomic nuclei” at JINR University Center for students of Moscow Physical-Technical

   Institute (September-December 2009, 2010).

3. V. O. Nesterenko: Lecture course “Introduction to physics of nanosystems”, Intern. Dubna University, (spring & autumn semesters (34+32 hours) 2009, 2010);

4. A. P. Severyukhin: exercises and exams, courses “General physics”, “Many-body system methods in nuclear physics”, Intern. Dubna University (2009, 2010).

5. R. V. Jolos: Lecture course ”Nuclear models”, 39 hours, Intern. Dubna University.

6. T.M. Schneidman: Lecture course "Selected topics in physics of nuclear structure and nuclear reactions"; JINR University Center.
7. R. G. Nazmitdinov: "Introduction to nanoscience and nanophotonics" (17 hours), I semester 2010; Intern. Dubna University
8. V. B. Belyaev: Lecture course "Nuclear Astrophysics" at Intern. Dubna University (2009, 2010).
9. E. A. Kolganova: Lecture course "Mathematical modeling and numerical methods" (February–June and September–December 2009, 2010).
10. V. S. Melezhik: Lecture course "General physics"(all the academic year), lecture course "History and methodology of physics" (September–December 2009, 2010), lecture course "Modern problems of quantum physics" (September–December 2009, 2010).
11. A. K. Motovilov: Lecture course and seminars on calculus for 1st year students (all the academic year 2009/2010), lectures and seminars on the course "Scattering theory for few-body systems" for 6th year students (September – December 2010).
12. A. I. Titov: Lecture course "Introduction to Particle and Nuclear physics", Institute of Laser Engineering, Osaka University, Japan, (August–September. 2010).