Momentum Dependence of the Ion-Ion Potential in a Microscopic Theory

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A momentum-dependent real potential for ion-ion scattering is calculated, from Reid's soft-core potential, with Brueckner's theory. The densities of the two nuclei are kept fixed and the Pauli principle is obeyed by rearrangement in momentum space. The calculated potential is increasingly attractive up to 2 fm^{-1} relative momentum per particle. At this relative momentum, the attraction does not decrease at small distances; for larger momenta, the potential finally becomes repulsive.

The scattering of heavy ions at relatively low energies above the Coulomb barrier is essentially a surface phenomenon. Elastic scattering is marked by strong absorption in the outer tail region of the real ion-ion potential which shields the elastic channel from regions of closer interpenetration.[1] Inelastic processes can be grouped roughly into two classes: (1) far grazing collisions which lead to direct reactions and few-nucleon transfer, and (2) deeply energy-relaxed reactions which can be successfully described as surface diffusion processes of the two touching nuclei.[2,3]

Though deeply inelastic collisions lead far away (in energy) from the entrance channel, they do not correspond to a strong interpenetration of the colliding nuclei. This can be understood in a semiclassical picture which is well suited for heavy-ion reactions: The relative motion and its coupling to a few collective degrees of freedom (e.g., neck formation or surface deformations) can be treated as a system of classical variables whose conservative forces are the gradients of the corresponding potential-energy surfaces while a coupling to the fast internal degrees of freedom leads to friction forces in this classical dynamical system.[4,5] The surface nature of heavy-ion reactions is then due to the strong repulsion exerted in the coordinate of separation after touching, due to the incompressibility of nuclear matter.

Potential energies as a function of separation of two nuclei have been calculated by several authors.[6-8] They adopt the limit of frozen densities and use an adiabatic approximation which gives the potential for small relative momenta of the scatterers. The potentials show strong repulsion for short distances.

One may ask whether this repulsion, and thus the surface nature of heavy-ion reactions, persist for larger relative momenta of the scatterers. Qualitatively, there are two effects working in opposite directions. The exclusion principle contributing to nuclear incompressiblity for nuclear matter at rest is weakened with increasing relative momentum. On the other hand, the nucleon-nucleon interaction looses its attraction for large momenta of the scattering nucleons, becoming repulsive at a relative momentum of about 2 fm^{-1}. From this, one may conjecture that an energy window exists for which the real part of the ion-ion potential becomes very weak, or even attractive.[9] If this is the case, internal excitations, or cooperative phenomena, of the interacting ions can be generated from areas of strong overlap, at energies which are not yet in the mesonic or relativistic domains.

In order to study these questions quantitatively we calculate the ion-ion potential starting from the Reid soft-core potential[10] and using Brueckner reaction-matrix techniques previously applied for nuclear matter.[11] Different from the previously considered nuclear matter system, we have here *two* infinite slabs of nuclear matter with a relative momentum K_r (Fig. 1). The first step of our calculation involves the energy of interaction between these slabs.

The calculation of the interaction potential involves two basic assumptions. (i) The frozen-density approximation. We do not allow the densities to relax during scattering. The exclusion principle has then to be satisfied by a rearrangement in momentum space. (ii) The local density approximation. This reduces the problem to a

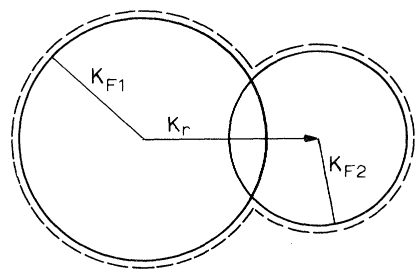

FIG. 1. The situation in momentum space for two pieces of nuclear matter with different Fermi momenta, k_{F1} and k_{F2}, at intermediate relative momentum K_r. The dashed contour shows the assumed rearrangement of occupied momentum states in order to avoid double occupancy in the overlapping area.

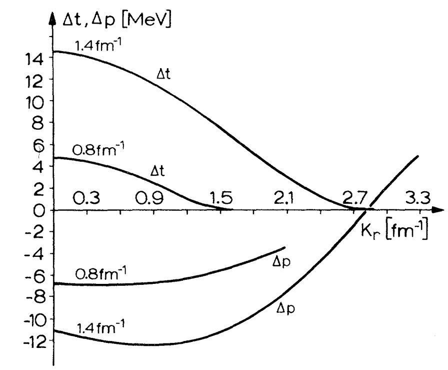

FIG. 3. The same as in Fig. 2, separated into the changes in kinetic, Δt, and potential, $\Delta \pi$, energies.

nuclear-matter calculation. Density gradient terms (like the Weizäcker correction) are neglected. Within this frame the interaction energy for two nuclei at a separation distance D and with relative momentum K_r is given by

$$V(D) = \int \epsilon(k_{F1}, k_{F2}, k_r)\rho(\vec{r}; D)\, d^3r, \tag{1}$$

where k_{F1}, k_{F2} refer to the local densities of the two scatterers and ρ is the total density. ϵ, the energy of interaction between infinite slabs, is the difference of the total energy per particle of the combined $(1+2)$ and the separated systems

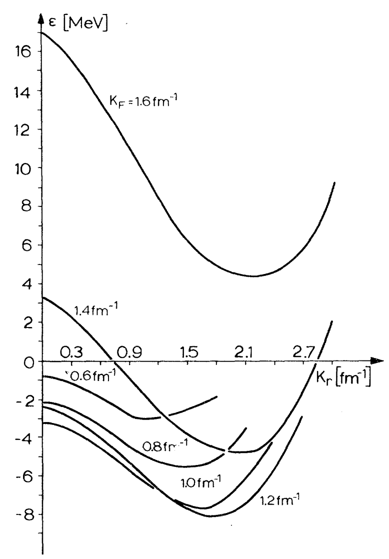

FIG. 2. The change in total energy per nucleon, ϵ, of two spatially overlapping pieces of nuclear matter with respect to the energies of the separated systems as a function of relative momentum per nucleon, K_r. The two pieces have equal densities, indicated by the corresponding Fermi momentum k_F.

(1 and 2)

$$\epsilon = \Delta\tau + \Delta\pi = \tau_{1+2} + \pi_{1+2} - (\tau_1 + \tau_2) - (\pi_1 + \pi_2). \tag{2}$$

The difference in kinetic energies, $\Delta\tau$, is due to a rearrangement in momentum space, induced by the Pauli exclusion principle (see Fig. 1).

The potential energy per particle, π, is obtained from a Brueckner G matrix:

$$\pi = \frac{2}{(2\pi)^3}\int_{\mathfrak{F}} \langle \vec{k}_1, \vec{k}_2 | G | \vec{k}_1, \vec{k}_2\rangle_A d^3k_1 d^3k_2 \Big/ \int_{\mathfrak{F}} d^3k, \tag{3}$$

where $\mathfrak{F}$ denotes the area of occupied states in momentum space (Fig. 1), and G is related to the bare nucleon-nucleon interacion v by

$$G = v + v\frac{Q}{e + i\eta}G \tag{4}$$

with Q, the Pauli operator, and e, the self-consistent energy denominator.

The bare interaction v is the Reid soft-core potential[10] in $^3S_1$, $^3D_1$, $^1S_0$, $^1D_2$, $^3D_2$, $^1P_1$, $^3P_0$, $^3P_1$, $^3P_2$, and $^3F_2$ states of relative motion. In general, G is non-Hermitian because of the boundary condition adopted in (4) which corresponds to outgoing waves for the real scattering states ocurring at finite momenta of the two scattering nuclei. Consequently, the calculated potential is an optical potential consisting of a real part and an imaginary part. In the scope of our calculation, the imaginary part is only important insofar as it renormalizes the real part.

In a first orientational calculation, we adopted the ordinary nuclear-matter self-consistent hole potential for $k_F = 1.6\ \text{fm}^{-1}$ in (4) and replaced the particle potential by zero and the boundary condition in (4) by the principal value.

Figure 2 gives total energy change per particle,

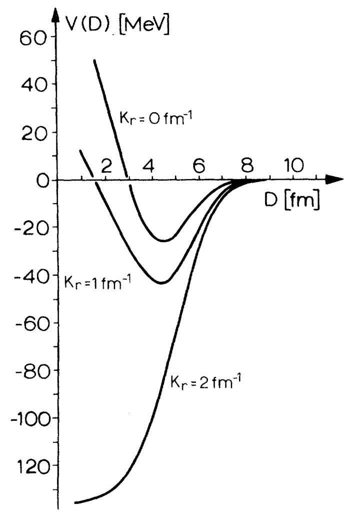

FIG. 4. ^{16}O-^{16}O potential as a function of relative separation for three values of the relative momentum K_r, for radius $R = 2.95$ fm and for surface diffuseness constant $a = 0.4$ fm. The total number of nucleon in one system is 35.03 (see text).

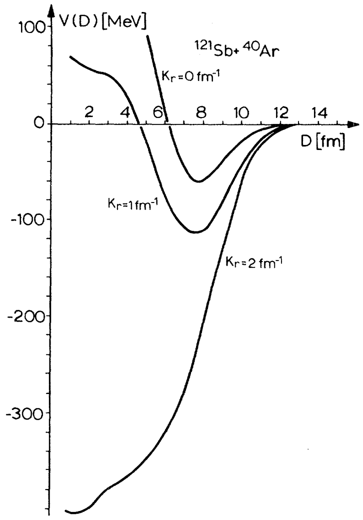

FIG. 6. The same as Fig. 4 for the system argon plus antimony.

ϵ, for two pieces of nuclear matter of equal densities overlapping completely in ordinary space, as a function of their relative momentum K_r. One observes that the expected window of a minimum in the interaction does occur for densities not too far away from saturation ($k_F \simeq 1.4\ fm^{-1}$) around $K_r \simeq 2\ fm^{-1}$.

Figure 3 shows the kinetic and potential terms of Eq. (2). One sees that the Pauli principle relaxes faster than the increase of repulsion in the interaction. This is the physical origin of the minimum in ϵ, shown in Fig. 2.

Figures 4-6 show some results for ion-ion potentials calculated according to Eq. (1). In these calculations the densities are of the Saxon-Woods form. The radii of the nuclei are chosen realistically but the central density is the saturating value for the Reid potential. Consequently the total nucleon number is larger than that of the real systems. Figures 4 and 5 show the influence of changing the surface diffuseness and it is pronounced for light systems such as ^{16}O. Figure 6 gives the potential for a heavy-ion system.

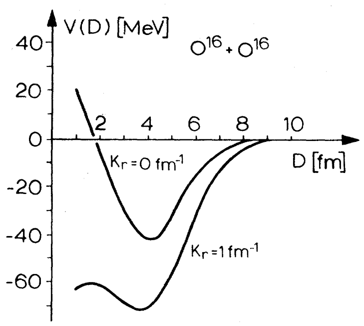

FIG. 5. The same as Fig. 4 for increased surface diffuseness, $a = 0.5$ fm.

We point out that the "pocket" increases with increasing momentum, and that the potential is attractive down to very small separation distances at the "window momentum" $K_r \simeq 2\ fm^{-1}$. The potentials for $K_r > 2\ fm^{-1}$ have not been calculated, but one can conclude from Fig. 2 that, for these momenta, the depth of the potential will decrease and eventually there will be repulsion.

Brink and Stancu[12] calculated a momentum-dependent ^{16}O-^{16}O potential using a Skyrme-interaction (SII) and found the most attractive potential at about 9 MeV/A while our results give us the higher value of about 20.7 MeV/A (c.m. energies). This difference is due to the unrelaistic, quadratic, momentum dependence for the Skyrme force.

The potentials shown cannot be used directly in a scattering calculation since, because of the momentum dependence, such a calculation has to be done self-consistently (presumably possible in a semiclassical local approximation). For quantitatively reliable results, the G-matrix equation, Eq. (4), has to be solved in a better approximation. We believe, however, that the general conclusion of a momentum (or energy) window at which heavy ions can easily interpenetrate remains valid.

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Search for Neutral-Weak-Current Effects in the Nucleus ^{18}F

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The circular polarization of the γ rays from the $1.08 \to 0.0$ MeV transition of ^{18}F has been measured to be $(-0.7\pm 2.0)\times 10^{-3}$, a value significantly smaller than predicted by recent calculations which include the effects of neutral weak currents.

The discovery of neutral-weak-current effects in reactions induced by high-energy neutrinos[1] has sharply increased interest in theories of neutral currents. One feature that may help distinguish between alternative theories[2] is the $\Delta T=1$ component of the parity-nonconserving (PNC) nucleon-nucleon weak forces. If neutral weak currents are PNC and isoscalar-isovector mixtures (as in the Weinberg-Salam model), $\Delta T=1$ parity mixing in nuclei may be enhanced by an order of magnitude over that predicted by the Cabibbo model[3] (charged weak currents only). All previously reported cases of nuclear parity mixing are either insensitive to the $\Delta T=1$ component of the weak force[4] or do not differentiate between $\Delta T=1$ and $\Delta T=0$ or 2 components.[5]

A favorable system for studying the $\Delta T=1$ PNC nucleon-nucleon force is provided by the 0^-, $T=0$, 1.08-MeV and 0^+, $T=1$, 1.04-MeV states of ^{18}F (see Fig. 1).[6] The circular polarization of the 1.08-MeV γ-ray transition from the 0^- state to the ground state directly measures the $\Delta T=1$ PNC matrix element between the 0^+, $T=1$ and 0^-, $T=0$ levels. Since the parity impurity in the 1.08-MeV state is well described by simple two-level mixing with the 1.04-MeV state, the circular polarization P_γ of the 1.08-MeV γ transition is given by

$$P_\gamma(1.08\text{ MeV}) = -2\langle 0^+, T=1|V^{\rm PNC}|0^-, T=0\rangle \times(\|M1\|/\|E1\|)\Delta E^{-1},$$

where $\Delta E = 39$ keV is the energy splitting between the two states, and $\|M1\|$ and $\|E1\|$ denote reduced matrix elements for the $M1$ and $E1$ transitions which de-excite the 1.04- and 1.08-MeV levels.

