TEMPERATURES FROM INTERMEDIATE ENERGY HEAVY ION COLLISIONS

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We propose a statistical model with complete equipartition among fragmentation and internal excitation degrees of freedom, but not with respect to the fragment's kinetic energy. The resulting internal temperature is very insensitive to the initial total energy in agreement with experiment.

Nucleus–nucleus collisions at intermediate energies provide a useful tool for studying nuclear matter properties far from the normal equilibrium state. However, despite the fact that a considerable amount of empirical data in the energy range from 30–200 MeV/nucleon already exists, the reaction mechanisms are far from understood. In fact, at present the theory is not able to describe the qualitative change in the reaction mechanism when going from low to higher energies (for recent review we refer to refs. [1,2]).

In this note we want to discuss a puzzling result concerning the two-particle correlation measurements at small relative momenta. These correlation experiments provide information about the space–time extent and the excitation energy density of the emitting source. From the relative population of the excited states of light fragments, such as Li and Be isotopes, a temperature of about 5 MeV has been extracted. This internal temperature is almost unaffected when changing both the projectiles and the beam energy in the range of 35–95 MeV per nucleon. For recent experiments we refer to refs. [3–5]. So, the remarkable constancy of this internal temperature suggest that there may be a limiting temperature of about 5–6 MeV above which fragment formation is strongly suppressed.

On has, however, to reconcile the relatively low internal temperature with another experimentally deduced quantity, the kinetic temperature, which is extracted by fitting the fragment's kinetic energy spectra by a moving source. This kinetic temperature turns out to be of the order of 10–25 MeV for bombarding energies from 35 to 94 MeV per nucleon.

In the present note we want to investigate this puzzling problem of having two distinct temperatures, by giving up the assumption that there is complete thermal equilibrium in the emitting source at the breakup stage. Especially, we would like to study the consequences of the assumption that there is equipartition with respect to the internal excitation energy of the fragments and to the different partitions of the fragmenting nucleus but there is no equipartition with the relative motion of the fragments. Thus, we assume that during the disassembly process those nucleons form fragments which are spatially correlated and still have the imprint of the kinetic energy distribution of the original hot nucleus.

In fact in any expansion, adiabatic or otherwise, the internal temperature of the fragments at breakup will be lower than the kinetic temperature. We illustrate this point briefly with the help of the coalescence model by calculating the ratio of fragments formed in its ground state to those formed in an excited state. The probability for forming a ground (W_0) or excited state (W_1), is simply given (cf. ref. [6]) by the overlap of the Wigner functions of the formed clusters with the single-particle distribution functions $f(r, p, t)$ characterizing the hot nuclear system.

For orientation we consider here the formation of

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a two-body cluster but the method may be generalized to include higher order clusters as well. The Wigner functions $\phi_0(1,2)$ and $\phi_1(1,2)$ of the two-body subsystem are calculated by using ground state and $1\hbar\omega$ excited state wave functions in a harmonic oscillator, the frequency of which is related to the size of the system via $\omega = 3\hbar/r_0^2 m$. The single particle distribution of the hot source has a form reminiscent of the solution of the Boltzmann equation

$$f(\boldsymbol{r},\boldsymbol{p},t) \sim \exp[-\boldsymbol{p}^2/2mT_0 - 3(\boldsymbol{r}-\boldsymbol{p}t/m)^2/2R_s^2] , \quad (1)$$

which describes the situation of a collisionless expanding system havig initially a temperature T_0 and a source radius R_s. The system described by eq. (1) has the property that the local temperature $T_{\rm local}$ of a comoving volume element decreases adiabatically with time so that $T_{\rm local}/n^{3/2} = \text{const.}$, where n is the particle density.

The ratio W_1/W_0 we relate to the internal temperature $T_{\rm int}$ as it is done in analyzing experimental data. We have

$$W_1/W_0 = \langle f(1)f(2)\phi_1(1,2)\rangle / \langle f(1)f(2)\phi_0(1,2)\rangle = \exp(-\Delta E/T_{\rm int}) . \quad (2)$$

Here ΔE stands for the energy difference of the ground and excited state of the two-body cluster considered.

The natural question arises under what condition the sudden formation of clusters as calculated by using the coalescence formula would lead to the envisaged Boltzmann occupancy factor in eq. (2). To check this we calculate the ratio W_1/W_0 for the special case that the particles of the source are confined in an oscillator potential of frequency ω_0 at a given temperature $T_{\rm local}$. The exact quantum distribution function reads

$$f(r,p) = [1/\cosh(\omega_0/2T_{\rm local})] \times \exp\{-\tanh(\omega_0/2T_{\rm local})(m\omega_0 r^2 + p^2/m\omega_0)\} , \quad (3)$$

and yields

$$W_1/W_0 = \{\cosh(\omega_0/T_{\rm local}) + \tfrac{1}{2}(\omega/\omega_0 + \omega_0/\omega)\sinh(\omega_0/T_{\rm local})\}^{-1} . \quad (4)$$

The desired Boltzmann occupancy factor emerges only in the case $\omega_0 = \omega$, i.e. when the hamiltonians governing the source and the fragments are the same.

For an extended source ($\omega_0 \to 0$) the ratio $W_1/W_0 = (1+\frac{1}{2}\omega/T_{\rm local})^{-1}$ emerges. This latter relation, which is also obtained when using the classical distribution (1) and assuming $R_s \gg r_0$, indicates that an occupancy factor emerges which is larger than the Boltzmann factor. Consequently, a higher temperature is deduced when expressing the population of nuclear states via eq. (2). This is illustrated in fig. 1 for two different initial temperatures $T_0 = 15$ MeV and $T_0 = 10$ MeV, respectively. It is seen that both the adiabatic temperature $T_{\rm local}$ and the internal one $T_{\rm int}$, which is calculated through eq. (2), fall off very steeply with decreasing density as soon as the expansion begins.

These results suggest that the temperature values extracted experimentally via the population of individual states of a specific fragment have to be associated with the temperature at which the systems break up into pieces and the fragments are formed.

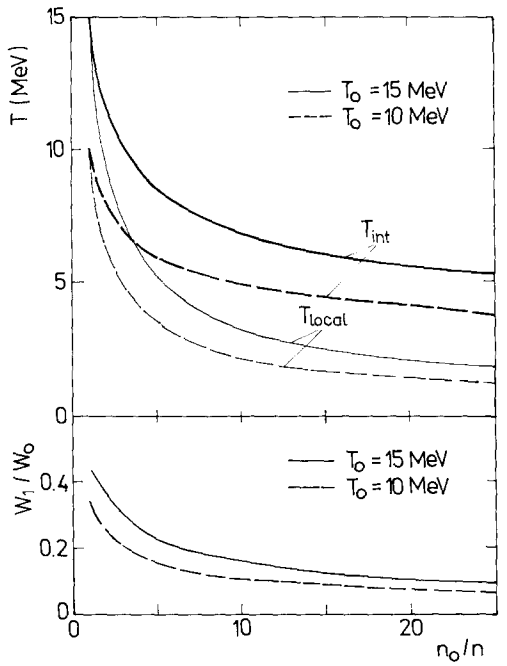

Fig. 1. Ratio (lower part) of the excited state and the ground state population of a fragment formed in an expanding source for two different initial temperatures T_0. In the upper part the adiabatic temperatures $T_{\rm local}$ (light lines) are compared to the internal ones $T_{\rm int}$ (heavy lines) which correspond to the occupation ratio W_1/W_0.

However, the energy spectra of the fragments, as well as the nucleons, are characterized by the higher kinetic temperature defining the transient stage at which free expansion begins. This reasoning seems to be in qualitative agreement with the experimental facts (see refs. [3–5]).

In the case of an isoentropic expansion, the breakup temperature is also of the order of about 5 MeV. However, the resulting fragment spectrum would not be a purely thermal one, but would have signatures of the collective flow velocity of the expanding source.

Now we briefly explain how one could incorporate this nonequilibrium behavior of the nuclear disassembly in a statistical description. As a starting model we use the statistical multifragmentation model (SMFM) [7], which in its spirit is similar to that of Koonin and Randrup [8], Fai and Randrup [9] and Gross and coworkers [10]. These fragmentation models can be viewed as an approach for the transient state of a hot nucleus at the breakup stage, whereby the possible states accessible for the disassembling system are determined by appropriate statistical weights given by the total entropy. Especially the SMFM [7] considers all possible partition channels in an effective free volume which is related to the total multiplicity and to a so-called crack width of $2d=2.8$ fm defined as the intersurface distance of neighboring fragments. The excitation energies and entropies of the hot fragments are calculated using a generalization of the liquid drop model to nonzero temperatures.

Usually in the statistical approaches the total entropy for the system is maximized at a given total excitation energy of the system. This implies that both the intrinsic states of the fragments and their translational energies are distributed according to a common temperature. But now we want to maximize the entropy under the constraint that the kinetic temperature T_{kin} characterizing the relative motion of the fragments is given. The internal temperature of the fragments at the breakup stage is then determined via the statistical decay modes considered in the SMFM.

We will investigate for the sake of illustration a system consisting of $A=40$ nucleons. In fig. 2 we show the internal temperature T_{int} of the fragments as a function of the kinetic temperature T_{kin} at given excitation energies per particle of the fragmenting nucleus. Most striking is that irrespective of the relatively large variations in the excitation energy per particle (4–12 MeV), the fragment temperature T_{int} remains almost flat as a function of the kinetic temperature T_{kin} and lies in the interval of 5–7 MeV. For comparison we also shown in fig. 2 the internal temperatures extracted experimentally via the population of particle unstable states of light fragments as a function of the kinetic temperatures determined by means of the moving source fits (see refs. [3–5]).

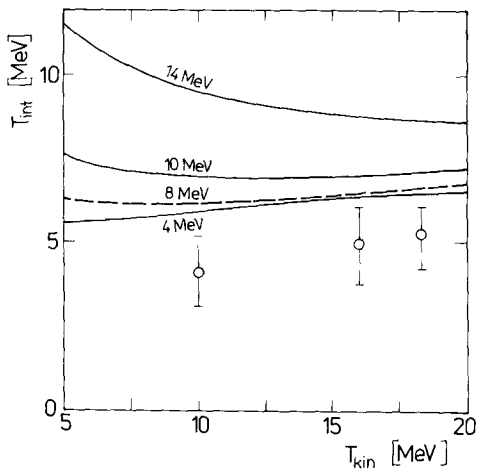

Fig. 2. Internal temperature T_{int} as a function of the kinetic temperature T_{kin} for different values (indicated) of the excitation energy per nucleon E^*/A. The data points represent temperatures T_{int} extracted from population ratios of ^{6}Li (ref. [4]) and T_{kin} from kinetic energy spectra averaged over different fragments produced in ion induced reactions on Au at 35 MeV [5], 60 MeV [3] and 94 MeV [4] bombarding energy per nucleon.

To understand better why the internal temperature does not change very much in the energy interval of interest when varying the kinetic temperature on a large scale, we display in fig. 3 the internal temperature T_{int} as a function of E^*/A for two given kinetic temperatures T_{kin}. One observes a plateau for $4<E^*/A<8$–10 MeV, which is connected with the fact that as soon as enough energy is stored into the nucleus to crack it, any additional energy goes mainly into surface energy in forming new fragments (see also ref. [7]). If enough fragments are formed the free-gas limit $T=2E^*/3A$ is approached. The slight backbend of these curves around $E^*/A\approx 4$ MeV is an outcome of SMFM [7] in which the effective free volume (see above) is a function of the fragment multiplicity that increases with E^*.

The results displayed in fig. 3 also indicate that

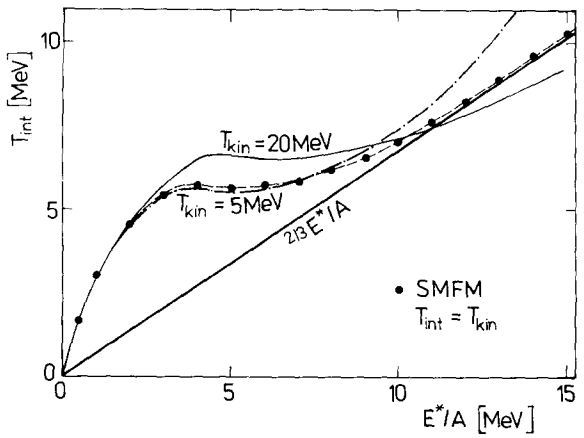

Fig. 3. The internal temperature calculated in the statistical model as a function of the excitation energy per nucleon E^*/A if the kinetic temperature T_{kin} is fixed to 5 and 20 MeV. The heavy dots are the results of the traditional multifragmentation model [7] assuming total equilibrium ($T_{int} = T_{kin}$).

when relaxing the assumption of a total equilibrium between the kinetic and internal temperature, the results are not much different from those where this assumption is maintained (heavy dots connected by dashed line). This is explained by the fact that for moderate excitation energies ($E^*/A \lesssim 8$ MeV) the kinetic energy of the fragments is significantly smaller than the energy deposited as internal energy. However, one has to keep in mind that when requiring total equilibrium the exponential slope of the fragment's energy spectra is determined by the internal temperature T_{int}, while now in our case the kinetic temperature T_{kin}, which has to be identified with the apparent temperature extracted by fitting the experimental spectra to Maxwell–Boltzmann distributions, determines the energy spectrum.

In summary, we have first studied a simple coalescence model to get a qualitative understanding of the experimentally observed fact that the temperature extracted via the population probability of certain states of light fragments turns out to be significantly smaller than the kinetic temperature characterizing the fragment's energy spectra. To describe the experimental data we proposed a modified statistical model which considers complete equipartition among fragmentation and internal excitation degrees of freedom but not with respect to the fragment's kinetic energy. We found that the internal temperature is of the order of 5–6 MeV for total excitation energies per particle in the interval 4–10 MeV and that it is not very sensitive with respect to variations of the kinetic temperature in the range of 5–20 MeV. These qualitative results are in agreement with data.

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References

[1] C.K. Gelbke and D.H. Boal, Prog. Part. Nucl. Phys. 19 (1987) 33.
[2] G.F. Bertsch and S. Das Gupta, Phys. Rep. 160 (1988) 189.
[3] J. Pochodzalla, C.K. Gelbke, W.G. Lynch, M. Maier, D. Ardounin, H. Delagrange, H. Doubre, C. Gregoire, A. Kyanowski, W. Mittig, A Peghaire, J. Peter, F. Saint-Laurent, B. Zwieglinski, G. Bizard, F. Lefebvres, B. Tamain, J. Quebert, Y.P. Viyogi, W.A. Friedman and D.H. Boal, Phys. Rev. C 35 (1987) 1695.
[4] Z. Chen, C.K. Gelbke, W.G. Gong, Y.D. Kim, W.G. Lynch, M.R. Maier, J. Pochodzalla, M.B. Tsang, F. Saint-Laurent, D. Ardounin, H. Delagrange, H. Doubre, J. Kasagi, A. Kyanowski, A. Peghaire, J. Peter, E. Rosato, G. Bizard, F. Lefebvres, B. Tamain, J. Quebert and Y.P. Viyogi, Phys. Rev. C 36 (1987) 2297.
[5] Z. Chen, C.K. Gelbke, J. Pochodzalla, C.B. Chitwood, D.J. Fields, W.G. Gong, W.G. Lynch and M.B. Tsang, Nucl. Phys. A 473 (1987) 564.
[6] G.F. Bertsch, Nucl. Phys. A 400 (1983) 221c.
[7] J.P. Bondorf, R. Donangelo, I.N. Mishustin, C.J. Pethick, H. Schulz and K. Sneppen, Nucl. Phys. A 443 (1986) 321; J.P. Bondorf, R. Donangelo, I.N. Mishustin and H. Schulz, Nucl. Phys. A 444 (1986) 460; H.W. Barz, J.P. Bondorf, R. Donangelo, I.N. Mishustin and H. Schulz, Nucl. Phys. A 448 (1986) 753; H.W. Barz, J.P. Bondorf and H. Schulz, Nucl. Phys. A 462 (1987) 742.
[8] S.E. Koonin and J. Randrup, Nucl. Phys. A 356 (1981) 223; A 474 (1987) 173.
[9] G. Fai and J. Randrup, Nucl. Phys. A 404 (1983) 551.
[10] B.-H. Sa and D.H.E. Gross, Nucl. Phys. A 437 (1985) 643; D.H.E. Gross, X.-Z. Zhang, and S.-Y. Xu, Phys. Rev. Lett. 56 (1986) 1544; X.-Z. Zhang, D.H.E. Gross, S.-Y. Xu and Y.M. Zheng, Nucl. Phys. A 461 (1987) 641; A 461 (1987) 668.