Microcanonical Thermodynamics, Fragmentation
“Phase-Transition”\(^a\), and the Topology of the N-body Phase Space

D.H.E. Gross

*Hahn Meitner Institut Berlin, Bereich TV,*
*and Fachbereich Physik der Freien Universität*
*Glienickerstr. 100, D14109 Berlin, Germany*

The general features of Microcanonical Thermodynamics (\(MT\)) as applied to the
fragmentation of hot nuclei and atomic clusters are discussed. \(MT\) is the most
fundamental form of any thermodynamics since Boltzmann. With modern computational techniques it is for the first time possible to explore it in nontrivial cases.
First and second order phase transitions in finite systems can unambiguously be
identified by the caloric equation of state \(T(E/N)\). The three characteristics of
phase transitions: The transition temperature, the latent heat and the \(\text{interphase surface tension}\) can be well defined and calculated in finite systems. Against common belief and in contrast to conventional (canonical) ther-

modynamics, \(MT\) of relatively small systems reflects the thermodynamic behavior of bulk systems in
great detail and surprising accuracy. By the Laplace transform from the micro-
canonical to the canonical ensemble many of the signatures of phase transitions
are smeared out or get lost. Also some of our believes about phase transitions
must be corrected: Microcanonical phase transitions show distinct peculiarities
very different from conventional ph-tr.: E.g. the specific heat is negative. Phase
separation is usually unsuitable to identify ph-tr. in small systems and in sharp
contrast to the caloric equation of state. In contrast to ordinary (canonical) ther-

dynamics \(MT\) allows the system to become \(inhomogeneous\) or to \(fragment\), at
first order phase transitions it gives insight into the coexistence region. Here the
form of the specific heat \(c(E/N)\) connects transitions of first and second order in
a natural way. The “phase transition” towards fragmentation is introduced. The
similarities and differences to the boiling of macro systems are pointed out.

Nuclear friction is likely responsible for the ergodic expansion of a multifragmented
nuclear source up to freeze-out densities of \(\sim 1/6\) normal density. Various recent
experimental results which seem to contradict our scenario are even confirming it
when investigated more thoroughly. The fragmentation “phase-transition” (multi-
fragmentation) in nuclei is by many reasons \(not\) the same as the liquid-gas transi-
tion in nuclear matter. The main reasons were already published 1984\(^a\).

1 Introduction

Our definition of Microcanonical Thermodynamics (\(MT\)) uses only mechanics.
The topology of the total accessible N-body phase space reflects (or implies?)
the behavior of many interacting N-body systems because their dynamics is

\(^a\)We use quotation marks to emphasize the limited analogy to standard phase transitions,
CRIS96,Catania May 27-31,1996. “Critical Phenomena and Collective Observables”
often chaotic. Then the dynamical evolution of many replica of the same system under identical macroscopic initial conditions follows the structure of the underlying N-body phase space. It is ergodic. In nuclear fragmentation this is presumably due to the strong and short ranged friction between moving nuclei in close proximity. Friction between atomic clusters is yet unknown but quite likely it exists there also.

First we have to discuss the concept of thermodynamics of small systems in general and especially of phase transitions. The first question is from which size on do phase transitions exist in small systems? How can one define them? Is a cluster of $\sim 100$ particles big enough? We will show that this is possible and one can unambiguously distinguish continuous (second order transitions) from discontinuous (first order) transitions by the form of the caloric equation of state $T_{thd}(E)$. Before proceeding further it is important to realize that isolated nuclei or clusters must be treated microcanonically. Usually there is no external heat- or particle bath which defines the temperature, the pressure, or the chemical potential. Microcanonical Thermodynamics is the proper theory for isolated small systems.

A microcanonical ensemble has some peculiarities: It does not have a positive definite heat capacity. In fact at a phase transition of first order the specific heat $c(\varepsilon)$ becomes negative in general. Therefore the classical signal of a peak in the specific heat is not useful to characterize a phase transition in small systems.

The second peculiarity of Microcanonical Thermodynamics is that it allows the system to become inhomogeneous: At first order phase transitions several regions of one phase coexist with other regions of the other phase. $MT$ allows differently to conventional thermodynamics the coexistence of regions with different energy-density. The partitional entropy of the spatial fluctuations is an important part of the total entropy. Small many-body systems and also large systems under long-range forces have an important new structural “phase transition” which does not exist in infinite homogeneous systems: They may fragment into few relatively large pieces. Typical examples are nuclear multifragmentation, as was predicted very early, but also the fragmentation of atomic clusters.

In this case the size fluctuations of the fragments at the transition are of the order of the size of the system itself. Then one cannot ignore the “droplets” compared to the nucleonic or monatomic vapor anymore. The “phase transition” is not determined alone by the equilibrium of the homogeneous liquid with the homogeneous gas, which in conventional (canonical) thermodynamics is controlled by the equality of the chemical potentials of liquid and gas as we are used to in conventional (grandcanonical thermodynamics). Often the
number of droplets is similar or larger than the number of free nucleons or monomers. This is one of the main lessons relevant for general physics which can be learned from nuclear fragmentations.

2 Microcanonical Thermodynamics

Microcanonical Thermodynamics explores the topology of the N-body phase space and determines how the volume $\Omega_N$ of the accessible phase space — more precisely the number of quantum states — depends on the fundamental globally conserved quantities of total energy $E = N \epsilon$, angular momentum $L$, mass (number of atoms) $N$, charge $Z$, linear momentum $p$, and last not least the available spatial volume $V$ of the system. This definition is the basic starting point of any thermodynamics since Boltzmann.

The entropy is defined as the logarithm of $\Omega$

$$S(E, V, N) = N s(\epsilon = E/N) = \ln(\Omega(E, V, N)) \tag{1}$$

and the thermodynamic temperature $T_{thd}$ is defined by

$$\beta = \frac{\partial S(E, V, N)}{\partial E} = \frac{\partial s(\epsilon)}{\partial \epsilon}, \tag{2}$$

$$T_{thd} = \frac{1}{\beta}. \tag{3}$$

By Laplace transform of $\Omega(E, V, N)$ one steps from the “extensive” variables like $E, V, N$ to the intensive ones like $T, P, \mu$. E.g. the Gibbs grand partition function and the Gibbs grand potential are then

$$Z(\beta, P, \mu) = \int_0^\infty \int_0^\infty \Omega(E, V, N) e^{-\beta(E + PV - \mu N)} \, dE \, dV \, dN, \tag{4}$$

$$G(\beta, P, \mu) = -T \ln[Z(\beta, P, \mu)]. \tag{5}$$

We take Boltzmann’s constant $k = 1$. In the same way one gets the canonical partition function and the free energy as

$$Z(\beta, P, N) = \int_0^\infty \int_0^\infty \Omega(E, V, N) e^{-\beta N(\epsilon + PV)} \, dE \, dV, \tag{6}$$

$$F(T, P, N) = -T \ln[Z(\beta, P, N)]. \tag{7}$$

3 Differences between microcanonical and canonical ensemble

According to van Hove a system of $N$ particles interacting via short range two-body attractive forces with hard cores is thermodynamically stable, the
thermodynamic limit of \( N, V \to \infty \left| _{N/V=\eta} \right. \) exists for such systems, intensive quantities like the specific energy have finite limiting values. Then, the thermodynamics derived from the microcanonical partition sum \( \Omega(E, N, V) \) and the one derived from the canonical \( Z(\beta, P, N) \) or the grand canonical partition function \( Z(\beta, P, \mu) \) (usually) coincide. Outside of phase transitions of first order, the relative fluctuations \( \Delta E/E \), or \( \Delta \eta/\eta \) vanish \( \propto 1/\sqrt{N} \).

This is quite different at phase transitions \((T = T_{tr})\) of first order and for finite systems where the microcanonical and the (grand)canonical ensemble differ essentially. In the canonical ensemble the energy fluctuations \( \Delta \varepsilon \) per particle remain finite even in the thermodynamic limit. \((\Delta \varepsilon)^2 |_{T_{tr}} \propto q_{lat}, \) the specific latent heat). Consequently, the difference between the microcanonical and the canonical ensemble persists at transitions of first order and we must expect both ensembles describe different physical situations.

Systems interacting via long range forces like unscreened Coulomb or the centrifugal force when they are rapidly rotating don’t have a thermodynamic limit and must be described by the microcanonical ensemble. Such systems fragment macroscopically into, in general several, regions of high density — condensed matter — and also into, in general several, regions of low density — vapor or may be empty space. Differently from conventional thermodynamics where systems which must be in a homogeneous phase at fixed temperature everywhere, here the system is most likely inhomogeneous. The inhomogeneities and their fluctuations are more important in characterizing the state of the system than any mean values. In contrast to thermodynamics of the homogeneous bulk, in small systems or large systems with long range forces the entropy connected to different partitions of the system is an important part of the total entropy. Familiar formulas like the one-particle entropy

\[
s_{sp} = - \sum_a n_a \ln(n_a)
\]  

are useless for calculating the total entropy.

4 First and second order phase transitions in small systems

Macroscopic systems have a discontinuity in the specific heat \( c_{bulk}(T) \) at first order phase transitions. \( c_{bulk}(T) \) may have a finite peak at \( T \approx T_{tr} \). On top of this there is a spike \( q_{lat} \delta(T - T_{tr}) \). With finite resolution it shows up as jump in \( c_{bulk}(T) \) by the latent heat \( q_{lat} \). A typical example is the specific heat of bulk sodium, fig. (1) at the melting transition. In contrast, a transition of second order is continuous at the transition temperature where \( c_{bulk}(T) \) has (in the example of the Ising model) a logarithmic singularity in \( T - T_{tr} \).
Microcanonical Thermodynamics gives new and deep insight into this. It will further allow to extend the concept of phase transitions to systems not treated before by thermodynamics like systems with long range forces or strongly rotating systems. We begin with the discussion of microcanonical phase transitions in standard model systems in which phase transitions of first and second order are well known. As example we take the 2-dimensional 10-states Potts model for which the asymptotic thermodynamics is even known analytically. In ref. we determined the three basic parameters of phase transitions of first order within Microcanonical Thermodynamics, the transition temperature $T_{tr}$, the specific latent heat $q_{lat}$, and the specific interphase surface entropy $\Delta s_{surf}$ for a system with nearest neighbor couplings. It was demonstrated that for surprisingly small systems the values of these three parameters are closer to their asymptotic values than in the canonical ensemble. This is so because most of the finite-size scaling is due to the large, but trivial, exponent in the Laplace transform eq.(6) from the micro- to the canonical partition sum.

Fig.(1): Specific heat of bulk sodium at atmospheric pressure from $7^\circ_C$ to $8^\circ_C$. The dashed line represents the specific heat calculated within the Debye model. The insert is a blow-up for $0 \leq T \leq 100K$.

The two types of phase transitions are distinguished by the form of the microcanonical caloric equation of state $T_{thd}(\varepsilon)$: A transition of first order has a backbending caloric equation of state $T^{-1} = \beta(\varepsilon)$ c.f. figure (2b). For a system with nearest neighbor interactions the area between $\beta(\varepsilon)$ and the “Maxwell” line $\beta = 1/T_{tr}$ is twice the interphase surface entropy $\Delta s_{surf}$. The left darkened area is the defect of entropy $\Delta s_{surf} = \int_{\varepsilon_1}^{\varepsilon_3} \beta d\varepsilon$ that the system ‘pays’ for introducing interphase surfaces, which it finally gets back when the whole system is converted to the new phase at $\varepsilon = \varepsilon_3$, right darkened area, and the interphase surface disappears. In the bulk the transition is discontinuous as function of $T$ or $\beta$ and “jumps” from the liquid ($\varepsilon \leq \varepsilon_1$) to the gas branch ($\varepsilon \geq \varepsilon_3$) of the caloric curve, fig.(2). As a function of the specific energy $\varepsilon$ the transition is however continuous. With rising $\varepsilon$ the system passes smoothly from the liquid phase over a mixed phase with coexisting large fluctuations of the two phases (“gas bubbles” and “liquid droplets”) to the pure gas phase.
when the specific energy is increased by the specific latent heat \( q_{\text{lat}} \). Figure (2a) shows the corresponding specific entropy \( s(\varepsilon) = \int_0^\varepsilon \beta(\varepsilon ') d\varepsilon ' \). The transition is characterized by the convex intruder in \( s(\varepsilon) \) of depth \( \Delta s_{\text{surf}} \). Figure (2c) shows the specific heat capacity

\[
c(\varepsilon) = \frac{\partial \varepsilon}{\partial < T >} = -\frac{\beta^2}{\partial \beta/\partial \varepsilon}
\]

as a function of the specific energy \( \varepsilon \). (Here numerical fluctuations in \( \beta(\varepsilon) \) in figure (2b) have been smoothed). One can see within the coexistence region of \( \varepsilon_1 \leq \varepsilon \leq \varepsilon_3 \) (shaded area in fig (2)), the microcanonical specific heat has two poles and becomes negative in between.

As the convex intruder in the specific entropy \( s(\varepsilon) \), fig.(2a), is forbidden by van Hove’s theorem in the canonical ensemble for an infinite number of particles, conventional thermodynamics of the bulk is blind in this energy interval. Here the Laplace transform eq.(6) has additional stationary points and the canonical bulk would be unstable. It can only see the branches of \( c(\varepsilon) \) in the regions \( \varepsilon \leq \varepsilon_1 \) and \( \varepsilon \geq \varepsilon_3 \). Thus the canonical specific heat \( c_{\text{bulk}}(T) \) will be positive and approach finite values at the transition of first order. At \( T = T_{\text{tr}} \), \( c_{\text{bulk}}(T) \) has an additional peak = \( q_{\text{lat}} \delta(T - T_{\text{tr}}) \).

If the specific latent heat \( q_{\text{lat}} \rightarrow 0 \) and the specific interphase surface entropy \( \Delta s_{\text{surf}} \rightarrow 0 \) the caloric equation of state gets only a saddle point at the transition. Then \( E(T)/N \) as well as \( T(E/N = \varepsilon) \) become single valued, the transition is continuous in the canonical as well as in the microcanonical ensemble. We have a phase transition of second order. The two poles of \( c(\varepsilon) \) merge and \( c(\varepsilon) \) or \( c_{\text{bulk}}(T) \) has a singularity at the transition point \( \varepsilon_{\text{tr}},T_{\text{tr}} \).

Both slopes of \( c(\varepsilon) \) are fully accessible in the canonical treatment of the heat capacity. Consequently, from the caloric equation of state \( T(\varepsilon) \) it is always possible to identify and distinguish both kinds of transitions. In Microcanonical Thermodynamics the relation between the two is very natural, transparent and simple.

It is further instructive that in finite realizations of the two-dimensional Potts model with \( q = 10 \) spin orientations at each lattice point it was not possible to see a clean separation into a compact region of ordered spins and a compact region of disordered spins at energies inside the coexistence region even for a lattice of 100 * 100 points. There were always several “gas bubbles” and “droplets” fluctuating over the lattice and prohibiting large interphase surfaces. Nevertheless the caloric equation of state \( T(\varepsilon) \) is already close to its asymptotic form. We can conclude from this observation that the other classical signal of a transition of first order, a clear separation of the two phases, is not a useful signal of a transition of first order in small systems.
5 Signals of a “phase transition” in nuclear fragmentation

In the review article ref.\textsuperscript{13} I proposed the caloric equation of state $T(\tilde{\varepsilon})$ for $^{131}\text{Xe}$. This had two anomalies compared to the standard parabolic dependence of a Fermi-gas $T \propto \sqrt{\varepsilon}$, fig.(13). One is quite pronounced with even a backbending of $T(\varepsilon^*)$ at $\varepsilon \sim 2.5\text{MeV}$, $T \sim 4.5\text{MeV}$ and a second one is less pronounced and has no backbending is at $\varepsilon \sim 5\text{MeV}$, $T \sim 6\text{MeV}$. The first one had a width $q_{\text{lat}} \sim 1\text{MeV/nucleon}$ and the second $\sim 1.5\text{MeV/nucleon}$ and may even be a transition of second order(?) \textsuperscript{?}. The latter one shall not be discussed here. This is likely the transition found recently at GSI\textsuperscript{14,15}. In the above review article the first anomaly was compared to the apparent slope temperature $T_{\text{app}}$ of evaporated $\alpha$-particles from $^{32}\text{S}+\text{Ag}$ and of $^{16}\text{O}+\text{Ag}$ from the Texas A&M group which show the same narrow anomaly at similar energies and temperatures.\textsuperscript{16,17} However, these data had too large error-bars to allow any firm conclusion. The interesting aspect of this first anomaly is that one may see it in theory as well as in experiment in $\alpha$-evaporation spectra even though in the model they are linked to the $\approx$ sudden opening of the IMF production.

Fig.(2) Specific entropy $s(\varepsilon) = \int_0^\varepsilon \beta_{\text{micro}}(\tilde{\varepsilon}) d\tilde{\varepsilon}$ vs. the specific energy $\varepsilon$ for the 2-dim. Potts model with $q = 10$ on a 200 $\times$ 200 lattice. In order to visualize the anomaly of the entropy the linear function $a + b\varepsilon$ ($a = s(\varepsilon = 0.17)$, $b = 1.42$) was subtracted. Because we use periodic boundary conditions one needs two cuts to separate the phases and the depth of the convex intruder is twice the surface-entropy.

b) Inverse temperature $\beta_{\text{micro}}(\varepsilon) = 1/T(\varepsilon)$ as directly calculated by MMMC
c) Specific heat $c(\varepsilon) = -\beta^2/(\partial^2 T/\partial \varepsilon)$. The canonical ensemble of the bulk jumps over the shaded region between the vertical lines at $\varepsilon_1$ and $\varepsilon_3$. This is the region of the coexistence of two phases one with ordered spins, the other with disordered spins. Here $c(\varepsilon)$ has two poles and in between it becomes negative. The canonical thermodynamics is blind to this region. Observe that the poles are inside $\varepsilon_1 \leq \varepsilon \leq \varepsilon_3$, i.e. the canonical specific heat remains finite and positive as it should.
Fig. (3) Experimental and theoretical (with MMMC) caloric equation of state, $T_{\text{app}}(\varepsilon^*)$ for $p$ and $\alpha$. The horizontal error-bars give the statistical uncertainty to extract the slope from the raw spectra in ref. 18. Different methods to determine the excitation energy lead essentially to a parallel up or down shift of the curves by the amount indicated by the vertical bars at the lowest and highest data point. The dash-dotted curve is a Fermi-gas calculation ($\varepsilon = T^2/8$).

The experiment by Chbihi et al. of incomplete fusion reactions of 701 MeV $^{28}$Si + $^{100}$Mo gives further evidence of this transition. We plot in fig. (3) $\varepsilon^*$, the excitation energy per nucleon vs. $T_{\text{app}}$, where $T_{\text{app}}$ is the slope of the raw evaporation spectra for protons and alpha particles. The curves give the $T_{\text{app}}(\varepsilon^*)$ dependence deduced from the microcanonical statistical multifragmentation model (MMMC) using its standard parameters. Also the experimental uncertainties for the proton and alpha curves are given. The similarity of the shapes of the experimental and simulated $T_{\text{app}}(\varepsilon^*)$ for the $\alpha$-spectra is quite evident. The differences between the shapes of these curves and the parabolic dependence (dotted curve) expected for a simple Fermi gas is clearly outside the experimental error margins indicating that some additional degrees of freedom, which are apparently included in the (MMMC) model, become significant in this energy range. The proton data are not so clear. They do indicate a similar anomaly but are more close to the parabolic Fermi-gas form than to the MMMC curve. Maybe the protons diffuse too fast out of the expanding soup of fragments and do not explore the structure of the accessible phase space sensitively enough.

The experimental analysis of the data provides the values of the mass $A$, charge $Z$, excitation energy $E^*$, and angular momentum $L$ of the source. The freeze-out radius $R_f$ was taken as its standard value of $2.2A^{1/3}$ fm, this means that we simulate a “phase transition” at constant volume. The results of MMMC calculations, performed by O. Schapiro, with these input values, were subjected to the same software filter as the experimental set-up which, most importantly, selects only those events with one big residue. The mass of the residue was chosen to be $A_{\text{res}} \geq 90$, which is close to $A_{\text{res}}$ estimated from
the experimental data (the mass of the residue could not be measured).

The theoretical value of $T_{\text{app}}$ was extracted from fitting, as was done for the raw experimental spectra. Similar to the experiment the individual temperatures for the different particle species are slightly different in MMMC from the thermodynamic temperature $T_{\text{thd}}$ c.f. figure (3). This interesting detail is due to the *intrinsic fluctuation* of $T_{\text{thd}}$ in the microcanonical ensemble — a result outside conventional thermodynamics.

The values of $R_f$ and $A_{\text{cs}}$ do not influence the general shape of the $T_{\text{app}}(\varepsilon^*)$ curves. However, the $T_{\text{app}}(\varepsilon^*)$ curves shift along the $T_{\text{app}}$ axis if different values of these parameters are used. The shifts produced by reasonable changes in $A_{\text{cs}}$ are larger than those produced by reasonable changes in $R_f$. We checked that the anomaly in $T(\varepsilon)$ is not due to the changes of the angular momentum from $L = 18.2$ to $48.8h$. It exists also at $L = 0$.

While the similarity of the shapes of the experimental and simulated $T_{\text{app}}(\varepsilon^*)$ is quite reasonable for $\alpha$, d, and t-particles, (see Fig.(3)) significant differences exist in detail. The simulated curves for d and t (not shown here) have the *same* S-shape as the data (and the $\alpha$-s) but are shifted towards lower values of $T_{\text{app}}$. The higher $T_{\text{app}}$ values of the experimental deuteron and triton spectra might again be an indication of a faster diffusion of these less bound fragments from a hotter stage of the early fragmented system. However, differently to the protons their diffusion is on the other hand slow enough to realize the transition in the underlying structure of the phase space. This interpretation is supported by a recent experiment by ref. showing a *simultaneous* early fragmentation and a smaller “decay” (diffusion) time for d and t.

The experimental data also suggest an association between the onset of IMF production and the anomaly in $T_{\text{app}}(\varepsilon^*)$. In both theory and experiment the multiplicity of Li-fragments just starts to rise at the same excitation. The success of the MMMC model in showing the same anomaly in the caloric equation of state in a nearly automatic manner has a potential far reaching consequence. The central assumption of this model is an ergodic mixing of the fragmented system filling the accessible N-fragment phase space with a well defined freeze-out region uniformly. This is quite different from sequential binary fragmentation, with *independent* motion of fragments outside the binary barrier. Here an anomaly in $E(T)$ would not have such a natural explanation.

Taking all evidence together: The anomaly in all four spectra (proton, deuteron, triton, and alpha) at the same excitation energy as predicted by MMMC and also the earlier data of the Texas A&M-group, in spite of their huge error bars there is in my opinion no doubt that this is a signal of a “phase transition” due to the relatively sudden opening of additional phase space at the onset of fragmentation, the lower one predicted in ref. Due to the strong
stochastic (dissipative) coupling of the various fragmentation channels this "phase transition" is felt in all channels even in the pure evaporation channels. The data of the Texas A&M group, the data of Chbihi et all, as well as the results of the Aladin collaboration are the first experimental caloric signals of a phase transition in a nucleus. The data of ref. have now smaller error bars and approach the higher (second order ?) transition mentioned above. They show a slow monotonic increase of $T$ with rising excitation. I.e. they do not anymore exhibit the dramatic plateau in $T(E/A)$ with the spectacular "specific latent heat" of $\sim 5$ MeV, close to the total binding energy per nucleon.

1. D.H.E. Gross. 1984 INS-RIKEN Symp., J.Phys.Soc.Jap. 54(1985)392
2. D.H.E. Gross and Meng Ta-chung. In 4th Nordic Meeting on Intermediate and High Energy Physics, page 29, Geilo Sportell, Norway, January 1981.
3. J.P. Bondorf. In C.H. Dasso, editor, Nuclear Physics, Proceedings of the Nuclear Physics Workshop, I.C.T.P., Trieste, Miramare, Italy, 5-30 October, 1981, pages 765–770, Amsterdam, New York, Oxford, 1982. North Holland.
4. D.H.E. Gross, L. Satpathy, Meng Ta-chung, and M. Satpathy. Zeit. Physik A 309:41, 1982.
5. Ludwig Boltzmann. Vorlesung über Gastheorie. Number 1. Akademische Druck-u. Verlagsanstalt, Graz, 1981.
6. L. van Hove. Physica, 15:951, 1949.
7. G. Borelius. Solid State Physics, 15:1, 1963.
8. R. Hultgren, R.L. Orr, P.D. Anderson, and K.K. Kelley. Selected Values of Thermodynamic Properties of Metals and Alloys. John Wiley and Sons, New York, 1963.
9. R.J. Baxter. J. Phys., C6:L445, 1973.
10. D.H.E. Gross, A. Ecker, and X.Z. Zhang. Annalen der Physik, in print, 1996.
11. A. Hülser. Z.Phys.B, 95:63–66, 1994.
12. M. Promberger. preprint, Erlangen, 1996.
13. D.H.E. Gross. Rep.Progr.Phys., 53:605–658, 1990.
14. J. Pochodzalla et al. Phys. Rev. Lett., 75:1040–1043, 1995.
15. C. Schwarz et al. GSI-Nachrichten, 01-96:5–16, 1996.
16. G. Nebbia et al. Phys.Lett. B, 176:20–25, 1986.
17. R. Wada et al. Phys.Rev.C, 39:497, 1989.
18. A. Chbihi et al. Phys.Rev.C, 43:652;43:666, 1991.
19. A. Chbihi, D.H.E. Gross, O. Schapiro, S. Salou, and L.G. Sobotka. submitted to Phys.Rev.Lett., 1995.
20. C.J. Gelderloos et al: Phys.Rev.Lett, 75:3082–3085, 1995.