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Significance of Temperature Measurements in Relativistic Nuclear Collisions*

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We discuss the importance of temperature measurements for probing the properties of dense, hot hadron matter in relativistic nuclear collisions. The effects of the collective matter flow are considered. It is pointed out that information about the existence of a limiting temperature $T^{\max} \gtrsim m_\pi$ can only be obtained from future experimental facilities with beam energies $E_{\mathrm{LAB}} > 5$ GeV/n. We also discuss the possibility of observing abnormal nuclear matter via a secondary, high temperature component in the particle spectra and via a shoulder in the pion multiplicity distributions.

The production of extended globs of strongly compressed and highly excited hadronic matter is one of the most intriguing aims of relativistic heavy ion physics. As up to now practically nothing is known about the properties of nuclear matter under these extremes, it is essential to elaborate which kind of observables can be used to extract information about the equation of state of the hot and dense matter formed in violent nuclear collisions.

In the present work we discuss the significance of the measurements of the particle spectra in nuclear collisions with high associated multiplicity for determining the temperature in the system experimentally. We show that information about the existence of a limiting temperature $T^{\max}$ in hadronic matter may be obtained by measurements of the bombarding energy dependence of the temperature, $T(E_{\mathrm{LAB}})$ at energies $E_{\mathrm{LAB}} > 5$ GeV/n. We discuss the effects of the collective matter flow on the measured temperatures. A detailed inspection of the large transverse momentum (p_T) regions of the spectra can provide further knowledge about the nuclear matter equation of state, e.g. about the possible existence of abnormal nuclear states.

A series of recent measurements at Berkeley and Dubna yielded interesting and promising results: Nagamiya and collaborators [1] measured the inclusive spectra of pions, protons and light composite fragments (d, t, He) with large p_T at $\theta_{\mathrm{CM}} = 90°$ for collisions of equal nuclei [C+C, Ne+NaF, Ar+KCl], and for the asymmetric reaction Ar+Pb at bombarding energies of $E_{\mathrm{LAB}} = 0.2$, 0.4, 0.8, and 2.1 GeV/n.

Quite analogous experiments have been carried out by the Moscow group in Dubna: Pb targets were bombarded with 3.6 GeV/n He and C projectiles from the Dubna Synchrophasotron and the inclusive proton spectra were measured [2]. In both experiments, the large p_T regions of the spectra are found to decrease exponentially with the fragments' kinetic energy in the equal speed system, i.e. in the nucleon-

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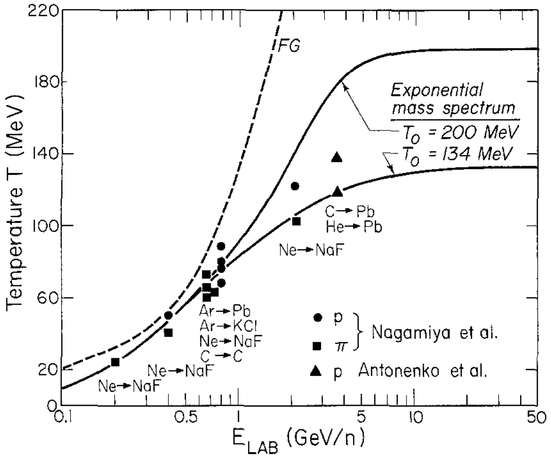


Fig. 1. Shows the experimentally determined slope factors $\tilde{T}(E_{\mathrm{LAB}})$. They are compared to the temperatures T calculated in the hydrodynamical model for a pure nucleon Fermigas (dashed) and for a classical ideal hadron gas with an exponentially increasing mass spectrum with $T_0=134$ MeV (lower solid curve) and $T_0=200$ MeV (upper solid curve), respectively

nucleon $(n-n)$ center of mass (c.m.) frame. The maxima of the invariant cross sections $\sigma_{\mathrm{inv}} \equiv E/p^2\, d^2\sigma/d\Omega dp$ for a given p_T, which are in general located at the rapidity of the $n-n$ c.m., can well be approximated by the expression

$$\sigma_{\mathrm{inv}}^{\max}(p_T)=\sigma_0 \exp(-E_{\mathrm{CM}}^*/\tilde{T})$$

with $E_{\mathrm{cm}}^*=(p_T^2+m^2)^{1/2}-m$, m being the rest mass of the particle detected. The slope factor $\tilde{T}$ - the "apparent temperature" [1] - of the inclusive proton and pion spectra is shown in Fig. 1 for various projectile-target combinations as a function of the bombarding energy E_{LAB}. Here we emphasize the following features of the data [1, 2]:

i) $\tilde{T}$ increases strongly with the bombarding energy. For the system Ne+NaF, the proton slope factors exhibit an increase from $\tilde{T}=49$, 75, to 122 MeV for $E_{\mathrm{LAB}}=0.4$, 0.8, 2.1 GeV/n, respectively [1].

ii) $\tilde{T}$ increases with the projectile and/or target mass at a given bombarding energy; for protons, $\tilde{T}=68$, 75 79, and 87 MeV for C+C, Ne+NaF, Ar+KCl, and the asymmetric system Ar+Pb at $E_{\mathrm{LAB}}=0.8$ GeV/n, respectively [1]. At $E_{\mathrm{LAB}}=3.6$ GeV/n, the $\tilde{T}$ values are 118 and 138 MeV for He and C projectiles, respectively [2]. This is the maximum value of $\tilde{T}$ observed up to now in relativistic nuclear collisions. Note that these $\tilde{T}$ values are close to the "asymptotic" slope factor observed in high energy proton-proton collisions $\tilde{T}_{pp} \approx 120$ MeV [3], which have been interpreted as evidence for a limiting temperature in hadron-hadron collisions [4].

iii) $\tilde{T}$ increases with the mass of the emitted fragment; for pions, $\tilde{T}$ is found to be systematically $\sim 20\%$ lower than in the corresponding proton data, which in turn are lower than the corresponding deuteron slope factors [1].

iv) $\tilde{T}$ increases with the associated fragment multiplicity $\langle M \rangle$, i.e., with decreasing impact parameters [1].

In the following we want to interpret these interesting experimental findings on the basis of a fluid dynamical approach to high energy nuclear collisions. First, let us compare the overall energy dependence (point *i*) of the observed "apparent temperatures" $\tilde{T}$ with the temperatures T calculated in the hydrodynamical model. This is done by solving the relativistic Rankine-Hugoniot equation [5, 6]

$$W^2-W_0^2+p(W/\rho-W_0/\rho)=0 \tag{2}$$

with an equation of state $W(\rho, T)$ appropriate for the hot dense hadronic matter. Here W, p, and ρ are the energy per baryon, pressure, and baryon number density in the rest frame of the compressed matter and $W_0=931$ MeV/n, $\rho_0=0.17\,\mathrm{fm}^{-3}$ are the ground-state energy and equilibrium density of nuclear matter, respectively. We use the following Ansatz to describe the density and temperature dependence of the equation of state:

$$W(\rho, T)=W_0+E_c(\rho)+E_T(\rho, T)+\Delta M(T) \tag{3}$$

The equation of state is decomposed into a zero temperature, density dependent "compression" part, $W_0+E_c(\rho)$, and a density and temperature dependent "thermal" part $E_T(\rho, T)+\Delta M(T)$, where E_T and ΔM describe the mean kinetic energy per baryon and mean mass increase due to resonance formation, respectively. For $E_c(\rho)$ we use the extended liquid drop model formula [5, 6]

$$E_c(\rho)=\frac{K}{18\rho\rho_0}(\rho-\rho_0)^2 \tag{4}$$

with the compression constant $K=200$ MeV. For the thermal part of the equation of state of the hadronic matter, let us consider a mixture of noninteracting hadrons in thermal equilibrium. The Fermi-Dirac and Bose-Einstein momentum space distributions for a relativistic ideal gas of particles of type i and of mass m_i, chemical potential μ_i, and statistical weight

$$g_i=(2S_i+1)(2I_i+1) \tag{5}$$

i.e., with spin S_i and isospin I_i, at a temperature T and occupying a volume V are ($\hbar=c=k_{\mathrm{Boltzmann}}=1$)

$$dN_i = \frac{g_i \cdot V p^2 dp}{(2\pi)^3 \{\exp[-(\varepsilon-\mu_i)/T] \pm 1\}} \tag{6}$$

where

$$\varepsilon = (p^2 + m_i^2)^{1/2}. \tag{7}$$

In the nonrelativistic classical (Boltzmann) statistics limit

$$dN_i = \frac{g_i V p^2 dp}{(2\pi)^3} \exp\left(\frac{\mu_i - m_i}{T}\right) \exp\left(\frac{-p^2}{2m_i T}\right) \tag{8}$$

The total number N_i and average energy density e_i (including rest mass) of particles of type i are obtained by integration of (6):

$$N_i = \frac{g_i V m_i^2 T}{2\pi^2} \times \sum_{n=1}^{\infty} \frac{(\mp)^{n+1}}{n} \exp\left(\frac{n\mu_i}{T}\right) K_2\left(\frac{nm_i}{T}\right) \tag{9}$$

and

$$e_i = \frac{g_i m_i^3 T}{2\pi^2} \times \sum_{n=1}^{\infty} \frac{(\mp)^{n+1}}{n}$$

$$\exp\left(\frac{n\mu_i}{T}\right) \left[K_1\left(\frac{n\mu_i}{T}\right) + \frac{3T}{n\mu_i} K_2\left(\frac{n\mu_i}{T}\right)\right] \tag{10}$$

This reduces to

$$N_i = g_i V \left(\frac{m_i T}{2\pi}\right)^{3/2} \exp\left(\frac{\mu_i - m_i}{T}\right) \tag{11}$$

and

$$e_i = \frac{N_i}{V}(m_i + \tfrac{3}{2} T) \tag{12}$$

in the nonrelativistic classical limit.

It has been argued [4] that one can account for the strong interactions in an ideal gas description of the hot, hadronic matter by using the complete hadronic mass spectrum, i.e., the complete density of states in phase space. However, the presently known hadrons have masses $\lesssim 2$ GeV and represent probably not the complete mass spectrum. In fact, Hagedorn's bootstrap hypothesis [4] ("A hadron is a fireball, which consists of fireballs, etc.") implies an infinite number of hadrons with an exponentially increasing hadronic mass spectrum, which in turn asks for a limiting temperature $T^{\max}$. Hamer and Frautschi [7] have generated, based on the bootstrap hypothesis, such an exponentially increasing mass spectrum using only the known hadrons as input. The density of states thus obtained is given by [8]

$$\rho^{HS}(m) = a\left(\frac{T_0}{m}\right)^3 \exp\left(\frac{m}{T_0}\right), \tag{13}$$

with $a = 0.56$/pion mass and $T^{\max} \equiv T_0 \equiv 134$ MeV to fit the experimentally known part of the mass spectrum between 1 and 2 GeV. Glendenning and Karant [8] have shown that in this model the production of antibaryons and strange hadrons is strongly suppressed in nucleus-nucleus collisions at energies $E^*_{\mathrm{CM}} \leqq 20$ GeV/n. This is due to the rather low achievable temperature $T < T_0 = 134$ MeV caused by an exponentially increasing mass spectrum. Also the number of mesons produced turns out to be very small due to the small available volume in the high compression stages of the collisions [8]. Let us for illustration study the pions, which, because of their low mass, give the by far largest contribution of all the mesons. For the intermediate temperature range, $1/3 < T/m_\pi < 3/2$, the number density ρ_π and energy density e_π of a relativistic ideal gas of pions can be approximated (to within 20% of the exact expressions) by the formulae [9]

$$\rho_\pi = 0.56\, \rho_0 (T/m_\pi)^4 \tag{14}$$

$$e_\pi = 1.85\, \rho_0 (T/m_\pi)^{9/2} m_\pi. \tag{15}$$

For $E_{\mathrm{LAB}} = 2$ GeV/n, where $T = 106$ MeV and $\rho = 7\rho_0$, the number of pions per baryon $\rho_\pi/\rho \approx 2.6\ 10^{-2}$, and the pion energy per baryon $e_\pi/e_B \approx 2.4\ 10^{-2}$. Therefore, the predominant contribution of the hadron gas to $W(\rho, T)$ is due to the excitation of the massive nonstrange baryonic resonances, N^*. Here we incorporate into the hadronic mass spectrum the nucleon, the Δ_{33} (1232 MeV) resonance, and an exponentially increasing resonance spectrum N^*, (13), which we use for $m \geqq 1.4$ GeV. For computational purposes, we discretize the exponentially increasing mass spectrum, so that for each mass interval of width m_π we consider one resonance with degeneracy $g_i = \rho^{HS}(m)$, (13).

Since the temperature achievable is limited to $T \lesssim T_0 \ll m_N c^2$, and on the other hand we are dealing with a nondegenerate Fermigas, $T \gg E_F$, the nonrelativistic classical expressions, (11) and (12) are appropriate to calculate the particle density ρ_i and energy density e_i of a given nucleon or hadronic resonance i:

$$\rho_i = \frac{N_i}{V} = \left(\frac{T}{2\pi}\right)^{3/2} g_i m_i^{3/2} \exp\left(\frac{\mu_i - m_i}{T}\right) \tag{16}$$

Since the heavy resonances N_i^* are excited states of the nucleon N, the rate equations in thermodynamic (chemical) equilibrium read

$$N + N \rightleftarrows N + N_i^* \tag{17}$$

from which the chemical potentials μ_i are obtained.

$$\mu_N + \mu_N = \mu_N + \mu_i, \tag{18}$$

i.e.

$$\mu_i = \mu_N \equiv \mu \tag{19}$$

Then, the total baryon number conservation ($\rho \equiv B/V$)

$$B = V \cdot \sum_i \rho_i \tag{20}$$

can be written as

$$B = V \left(\frac{T}{2\pi}\right)^{3/2} e^{\mu/T} \sum_{i=1}^{\infty} g_i m_i^{3/2} \exp(-m_i/T) \tag{21}$$

Therefore, the probability of forming a particle i at temperature T is given by

$$\lambda_i \equiv \frac{\rho_i}{\rho} = \frac{g_i m_i^{3/2} \cdot \exp(-m_i/T)}{\sum_{k=1}^{\infty} g_k m_k^{3/2} \exp(-m_k/T)} \tag{22}$$

or

$$\lambda_i = \frac{\tau_i \exp(-E_i/T)}{\sum_{k=1}^{\infty} \tau_k \exp(-E_k/T)} \tag{23}$$

with the effective degeneracy

$$\tau_i = \frac{g_i}{g_N} \left(\frac{m_i}{m_N}\right)^{3/2} \tag{24}$$

and the effective excitation energy

$$E_i = m_i - m_N \tag{25}$$

where $g_N = 4$ and $m_N = 939$ MeV are the statistical weight and the mass of the nucleon, respectively. The baryon number conservation can also be written as

$$\sum_{i=1}^{\infty} \lambda_i = 1 \tag{26}$$

The thermal energy per baryon is given as the sum of the kinetic energy of the resonances

$$E_T = \sum_{i=1}^{\infty} \lambda_i \cdot \tfrac{3}{2} T = \tfrac{3}{2} T \tag{27}$$

and the average increase of the rest mass due to the excitation of heavy resonances

$$\Delta M = \sum_{i=1}^{\infty} \lambda_i E_i \tag{28}$$

The total pressure of the system

$$P = P_C + P_T \tag{29}$$

is the sum of the compressional (zero temperature) pressure

$$P_C = \frac{K}{18}(\rho^2 - \rho_0^2) \tag{30}$$

and the thermal pressure P_T, which in turn is the sum of the partial pressures of the various resonances,

$$P_T = \sum_{i=1}^{\infty} P_{T_i} = \sum_{i=1}^{\infty} \rho_i T = \rho T \tag{31}$$

The entropy per particle i is given by

$$S_i = \tfrac{5}{2} + \ln \left[\frac{g_i}{\rho_i} \left(\frac{m_i T}{2\pi}\right)^{3/2}\right] \tag{32}$$

and, consequently, the entropy per baryon is

$$\begin{aligned} S &= \sum_{i=1}^{\infty} \lambda_i S_i \\ &= \tfrac{5}{2} + \ln \left[\frac{g_N}{\rho} \left(\frac{m_N T}{2\pi}\right)^{3/2}\right] + \sum_{i=1}^{\infty} \lambda_i \ln \left(\frac{\tau_i}{\lambda_i}\right) \\ &= \tfrac{5}{2} + \ln \left[\frac{g_N}{\rho} \left(\frac{m_N T}{2\pi}\right)^{3/2}\right] + \left(\sum_{k=1}^{\infty} \tau_k \exp \frac{-E_k}{T}\right)^{-1} \\ &\quad \cdot \sum_{i=1}^{\infty} \tau_i \exp \frac{-E_i}{T} \left[\frac{E_i}{T} + \ln \sum_{k=1}^{\infty} \tau_k \exp \frac{-E_k}{T}\right] \end{aligned} \tag{33}$$

Figure 1 shows the comparison of the experimentally determined apparent temperatures $\tilde{T}$ with the temperatures T as calculated in the hydrodynamical model, obtained by solving (2). Three different thermal parts of the equation of state have been investigated: *a)* the pure nucleon gas, i.e. the resonance formation is not included, *b)* an exponentially increasing mass spectrum, (13), with $a = 0.56$/pion mass and $T_0 = 134$ MeV, and *c)* an exponentially increasing mass spectrum with $T_0 = 200$ MeV. This higher limiting temperature is suggested by recent calculations on the critical temperature for the hadron-quark matter phase transition [10].

The pure nucleon gas seems to be ruled out experimentally since the predicted temperatures are much larger than the experimentally determined "apparent temperatures" for $E_{\mathrm{LAB}} \gtrsim 1$ GeV/n. Resonance formation must be included to decrease the achievable temperatures in order to resemble the apparent trend in the bombarding energy dependence of $\tilde{T}$. On the other hand, the highest bombarding energy presently available, $E_{\mathrm{LAB}} = 3.6$ GeV/n, seems to be too low to determine the asymptotic form of the hadronic mass spectrum. To determine whether a plateau in $T(E_{\mathrm{LAB}})$ actually exists is an important project for future experimental facilities. Higher bombarding energies, $E_{\mathrm{LAB}}>5$ GeV and heavier projectiles, e.g. Pb, are necessary to answer this important question finally. Only at these high energies could a possible exponentially increasing hadronic mass spectrum and a limiting temperature T_0 manifest themselves.

However, rather than the trend for a plateau in $\tilde{T}(E_{\mathrm{LAB}})$ at high energies, the agreement of the absolute values of the measured "apparent temperatures" $\tilde{T}$ with the calculated temperatures should be taken cautiously. $\tilde{T}$ overestimates the actual temperature by about 10 to 20 percent, since $\tilde{T}$ is determined from the invariant cross sections

$$\sigma_{\mathrm{inv}}=\sigma\exp(-E^*_{\mathrm{CM}}/\tilde{T})=E\cdot d^3\sigma/dp^3$$
$$=(m^2+p^2)^{1/2}\exp(-E^*_{\mathrm{CM}}/T) \tag{34}$$

Furthermore, for the lighter systems investigated (C + C, Ne + NaF) single nucleon-nucleon scattering can contribute considerably to the spectra [11] and hence contaminate the thermal contribution. In a simple thermal model $\tilde{T}$ should stay constant at a given bombarding energy for symmetric systems and drop for asymmetric systems, respectively. In contrast, the data exhibit a systematic increase of $\tilde{T}$ with the total number of particles in the system, see point *(ii)* above. In particular, $\tilde{T}$ increases when going from the symmetric system, Ar + KCl, to the asymmetric system, Ar + Pb, opposite to what one would naively expect. This mass dependence of $\tilde{T}$ indicates enhanced equilibration in the case of heavy mass nuclei, possibly due to increasing collective interactions [12].

Another interesting feature of the data is the ~20% difference in the $\tilde{T}$ values for protons and pions [1], see point *iii)*. Siemens and Rasmussen [13] argued that this difference may reflect the isentropic hydrodynamical flow of the expanding matter.

In thermal equilibrium the thermal velocities of the nucleons are much smaller than those of the pions. An additional collective flow velocity will therefore shift the proton spectra to larger transverse momenta, while the pion spectra are less affected. A mean flow velocity β and a mean temperature have been fitted [13] to the Ar (800 MeV/n) + KCl $\rightarrow p, \pi$ data at $\theta_{\mathrm{CM}}=90°$, using

$$\frac{d^3\sigma}{dp^3}=\tilde{N}\exp(-\gamma E/T)\left[\left(\gamma+\frac{T}{E}\right)\frac{\sinh\alpha}{\alpha}-\frac{T}{E}\cosh\alpha\right] \tag{35}$$

where $\gamma=(1-\beta^2)^{-1/2}$, $\alpha=\gamma\beta p/T$ and $\tilde{N}$ gives an absolute normalization. Fluid dynamical calculations of the expansion of a compressed fireball [14] demonstrate that the flow velocity $\mathbf{v}(\mathbf{r})$ depends linearly on the distance $\mathbf{r}$ from the center of mass of the fireball, $\mathbf{v}(\mathbf{r})\simeq\mathrm{const.}\cdot\mathbf{r}$, and that in the break-up stage the densities and temperatures in the fireball are nearly independent of $\mathbf{r}$. Using (35) for a finite number of fluid cells, the particle spectra can be calculated for an isentropic expansion process as a function of the break-up density ρ_{BU}.

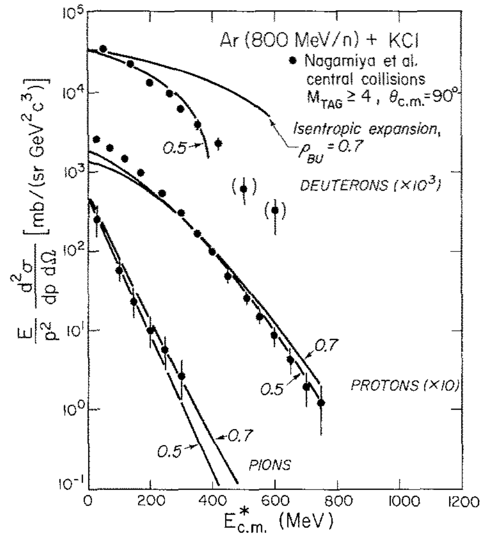


Fig. 2. Depicts the spectra of π^-, p, and d emitted from high multiplicity selected collisions of Ar (800 MeV/n) + KCl [1]. The data [1] are compared to the hydrodynamic calculations described in the text for two different break-up densities $\rho_{\mathrm{BU}}/\rho_0=0.5$ and 0.7

Figure 2 shows the calculated spectra of protons, pions, and deuterons for the reaction Ar (800 MeV/n) + KCl at two different break-up densities $\rho_{\mathrm{BU}}=0.5$ and $0.7\,\rho_0$. The calculations are compared to the high multiplicity selected, i.e., central triggered, data of Nagamiya et al. [1]. The absolute normalization is adjusted to the data. We want to point out that the shape of the proton and pion spectra are apparently not very sensitive to the exact value of the break-up density. Total energy conservation constrains the collective flow and the random thermal motion in such a way that the "apparent temperature" stays almost constant during the isentropic expansion. In fact, the slope in the high energy parts of the p and π spectra are reasonably well described by the initial temperature as shown in Fig. 1.

On the other hand, the deuteron spectra are strongly affected by the expansion. This is due to the large mass of the deuterons: v_{thermal} is always much smaller than v_{flow}. This may be used to determine the freeze-out density from experiment. Hence, although the temperature in the break-up moment is much smaller than the initial temperature, the observed proton and pion spectra can still carry the information (with $\approx 30\%$ uncertainty in the absolute value) about the initial temperature.

Another important question connected with the temperature measurements is whether from these data information can be obtained about abnormal nuclear matter, e.g., density isomers [15, 16], pion condensates [12, 16], or quark matter [6], possibly formed in relativistic nuclear collisions.

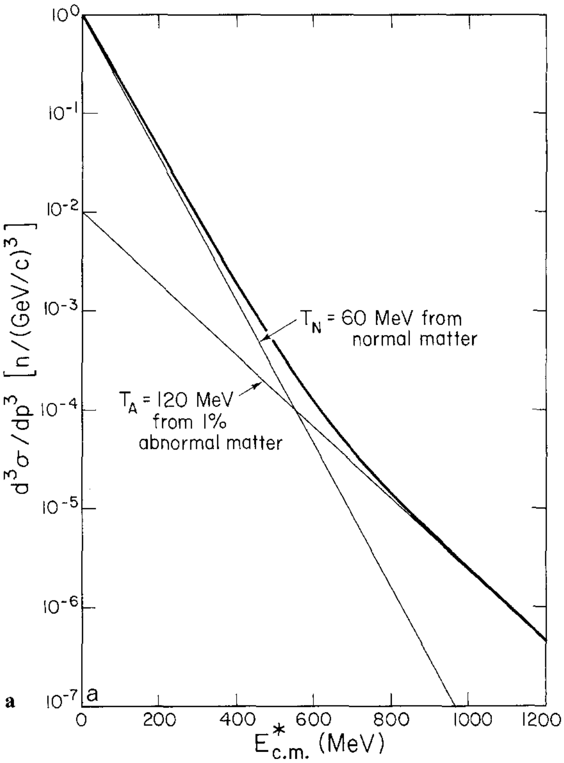
It has been predicted theoretically [6, 17] that the formation of abnormal nuclear matter (density isomers) in central nuclear collisions should be experimentally observable by a jump in the temperature and particle (e.g. π, K, α, or n) production excitation functions. This happens just above the bombarding energy for which the critical density ρ_c between normal and abnormal state is reached. The sudden release of the condensation energy leads to a strong entropy production [6, 17], thus also increasing the temperature of the matter.

However, more detailed three dimensional fluid dynamical calculations [18] show that the amount of abnormal matter formed drops drastically when the impact parameter is increased. For inclusive, i.e. impact parameters averaged data the threshold increase of the temperature, entropy and particle production rate may therefore be difficult to observe. The spectra will be dominated by the contributions of the normal matter, as illustrated in Fig. 3.

However, even small amounts of very hot abnormal matter may be detected in inclusive measurements, since the spectra would reveal such a second component via a second, larger slope constant in the large transverse momentum tails (see Fig. 3a). Only if the temperature difference between the normal and the abnormal state is not very large, e.g. on the order of $<20\%$ as shown in Fig. 3b, it will be hard to detect even a large ($\sim 10\%$) contribution from an abnormal state. The inclusive spectrum can then exhibit the higher second slope constant only at very large transverse momenta, i.e. at very small absolute cross sections.

The measurement of the particle spectra for central, i.e., high multiplicity selected collisions, can offer a way out of this dilemma because of the increased probability of producing an abnormal state. But not only central collisions offer a chance of producing and detecting abnormal matter. The collision processes are not deterministic, i.e., a given initial configuration can yield different results. For example, even a central collision at a bombarding energy exceeding the critical energy necessary to overcome the barrier between the normal and (hypothetical) abnormal state will have a finite probability of producing normal matter only. If this probability is large, the same problems as for the inclusive

a



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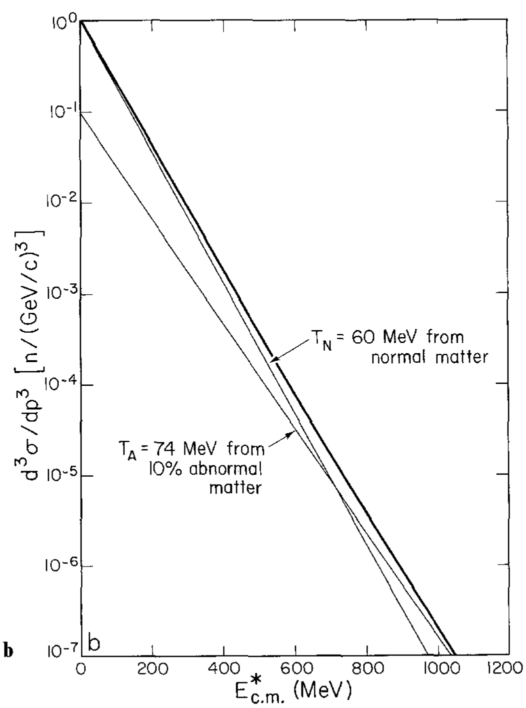

Fig. 3. The influence of a small amount of abnormal nuclear matter on the spectra of the emitted particles is shown: a second, high temperature component T_A is to be expected, which should be clearly observable if T_N and T_A differ sufficiently (Fig. 3a) but will be harder to detect otherwise (Fig. 3b).

measurements would persist for the central data. Vice versa, collisions at large impact parameters and/or bombarding energies below the critical energy can have a finite probability for a phase transition into an abnormal state. The peripheral collisions can lead to the formation of rather cold abnormal matter, which would then move with velocities close to the incident projectile and target velocities. Due to the condensation energy, their temperatures will be still larger than the ~8 MeV observed for the normal energy independent fragmentation/evaporation processes in peripheral collisions. Thus, also in peripheral collisions these states may be detected experimentally by secondary slopes in the spectra. Indeed, experimental evidence for such a second temperature has been reported recently [19] for periperal collisions. It is important to establish conclusively that this second component is not due to a contamination of the peripheral collisions by the hydrodynamic bounce-off [20] events with their large transverse momentum contribution.

Because of the difficulties in separating two similar temperature components from the inclusive spectra, see Fig. 3b, we have looked for a different, more sensitive observable, which can provide information on a high temperature component. It has been pointed out [6, 17, 21] that a threshold increase in the mean pion multiplicity excitation function would result from central collisions when most of the compressed matter undergoes the phase transition into an abnormal state. However, as discussed above (see Fig. 3b) the abnormal matter may be formed in a small fraction (e.g. on the order of a few percent) of the collisions only.

Then the pion production excitation function, i.e., the mean number of created pions versus the bombarding energy will exhibit a smooth linear behavior [22] and not a dramatic threshold increase. We will show now that the pion multiplicity distribution $P(n_\pi)$ can provide signatures of abnormal matter even under such unfavorable conditions. The pion multiplicity distribution is given by a Poisson distribution [23].

$$P(n) = e^{-\bar{n}}\,\bar{n}^n/n! \tag{36}$$

where $\bar{n}$ is the average number of pions, which in a thermal model depends on the temperature and volume of the pion emitting source. Using a relativistic fluid dynamical model with viscosity effects included, $\bar{n}_\pi(E_{\rm LAB})$ has been calculated [22] using a conventional nuclear equation of state. The calculations agree with the available data [24] for the rather light system Ar+KCl over a wide range of bombarding energies.

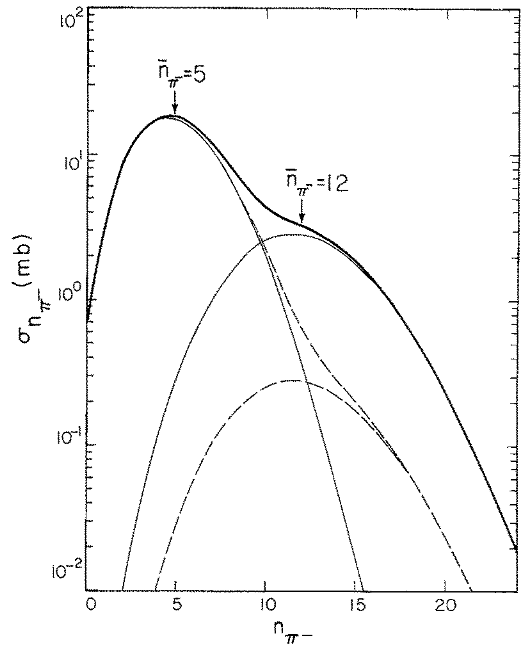

Fig. 4. The pion multiplicity distribution is shown for central collisions of $A_P = A_T = 100$ at $E_{\rm LAB} = 0.8$ GeV/n. Abnormal nuclear matter, if also produced, would lead to the shoulder at high multiplicities

In Fig. 4 we show the multiplicity distribution $P(n_\pi)$ of negative pions resulting from a central collision of $A_P = A_T = 100$ at $E_{\rm LAB} = 0.8$ GeV/n and associated with the same scenario as discussed in Fig. 3b, namely a normal temperature $T_N = 60$ MeV and a 10% contamination from abnormal matter $T_A = 74$ MeV. This ~20% difference in T leads to a quite dramatic difference in $P(n_\pi)$. The higher temperature associated with an abnormal state results in a distinct shoulder in $P(n_\pi)$. Even an only 1% contamination of abnormal matter should be experimentally observable at high multiplicities. Hence, measurements of the pion multiplicity distribution in central collisions of heavy systems can provide signatures for abnormal matter formed in only a small fraction of the reactions.

In conclusion, we have discussed the importance of the temperature measurements at future experimental facilities with higher beam energy $E_{\rm LAB} > 5$ GeV/n and higher projectile masses to search for a limiting temperature $T^{\max}$ in hot hadronic matter. We have considered the effects of the collective matter flow on the measured temperatures. We have also discussed the possibility of observing abnormal nuclear matter formed in central and peripheral collisions of heavy nuclei via a secondary, high temperature component in the spectra of emitted particles and via a shoulder in the multiplicity distribution $P(n_\pi)$ of pions created in such collisions. The experiments proposed above will provide a direct test of the equation of state of dense, hot hadronic matter in the near future.

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References

1. Nagamiya, S., Lemaire, M.C., Moeller, E., Schnetzer, S., Shapiro, G., Steiner, H., Tanihata, I.: LBL preprint 12123 (1981) and references therein
2. Antonenko, V.G., Galitskij, V.M., Grigor'yan, Yu.I., Ippolitov, M.S., Karadjev, K.V., Kuz'min, E.A., Manko, V.I., Ogloblin, A.A., Paramanov, V.V., Tsvetkov, A.A., Vinogradov, A.A.: Pis'ma JETP **28**, 609 (1978) and Moscow preprint IAE-310 (1979)
3. Laasanen, A.T., Ezell, C., Gutay, L., Schreiner, N.W., Schubelin, P., von Linder, L., Turkot, F.: Phys. Rev. Lett. **38**, 1 (1977)
4. Hagedorn, R.: Suppl. Nuovo Cim. **III.2**, 147 (1965) and In: Cargese Lectures in Physics. E. Schatzman (ed.), Vol. VI. 1973
5. Scheid, W., Müller, H., Greiner, W.: Phys. Rev. Lett. **32**, 741 (1974)
Hofmann, J., Stöcker, H., Heinz, U., Scheid, W., Greiner, W.: Phys. Rev. Lett. **36**, 88 (1976)
6. Stöcker, H., Greiner, W., Scheid, W.: Z. Phys. A – Atoms and Nuclei **286**, 121 (1978)
Stöcker, H., Graebner, G., Maruhn, J.A., Greiner, W.: Phys. Lett. **95**B, 192 (1980)
7. Hamer, C.J., Frautschi, S.C.: Phys. Rev. D4, 2125 (1971)
8. Glendenning, N.K., Karant, Y.: Phys. Rev. Lett. **40**, 374 (1978) and Phys. Rev. C**21**, 1501 (1980)
9. Sobel, M., Siemens, P.J., Bondorf, J.P., Bethe, H.A.: Nucl. Phys. A**251**, 502 (1975)
10. Satz, H.: Invited talk at the 5th High Energy Heavy Ion Summer Study, Berkeley, May 1981
11. Schürmann, B., Chemtob, M.: Z. Phys. A – Atoms and Nuclei **294**, 371 (1980) and Nucl. Phys. A**336**, 508 (1980)
12. Ruck, V., Gyulassy, M., Greiner, W.: Z. Phys. A – Atoms and Nuclei **277**, 391 (1976)
Gyulassy, M., Greiner, W.: Ann. Phys. **109**, 485 (1977)
13. Siemens, P.J., Rasmussen, J.O.: Phys. Rev. Lett. **42**, 880 (1979)
14. Csernai, L.P., Barz, H.W.: Z. Phys. A – Atoms and Nuclei **296**, 173 (1980)
15. Feenberg, E., Primakoff, H.: Phys. Rev. **70**, 980 (1946)
Bodmer, A.R.: Phys. Rev. D4, 1601 (1971)
Migdal, A.B.: Zh. Eksp. Teor. Fiz. **63**, 1993 (1972)
Ne'eman, Y.: In: Physics of Dense Matter. I.A.U. Symposium, Boulder, Colorado, 1972
Lee, T.D., Wick, G.C.: Phys. Rev. D**9**, 2291 (1974)
Boguta, J.: LBL preprint 12333 (1981)
16. Migdal, A.B.: Rev. Mod. Phys. **50**, 107 (1978) and references therein
Weise, W., Brown, G.E.: Phys. Rep. **27**, 1 (1976)
17. Stöcker, H., Maruhn, J.A., Greiner, W.: Phys. Lett. **81**B, 303 (1979)
18. Stöcker, H., Maruhn, J.A., Greiner, W.: Z. Phys. A – Atoms and Nuclei **290**, 297 (1979)
Z. Phys. A – Atoms and Nuclei **294**, 125 (1980), Prog. Part. Nucl. Phys. **4**, 133 (1980)
19. Baumgardt, H.G., Friedländer, E., Schopper, E., Otterlund, I.: Contributions to the 5th High Energy Heavy Ion Summer Study, Berkeley, May 1981
20. Stöcker, H., Maruhn, J.A., Greiner, W.: Phys. Rev. Lett. **44**, 725 (1980); Z. Phys. A – Atoms and Nuclei **293**, 173 (1979); Prog. Part. Nucl. Phys. **4**, 133 (1980); Stöcker, H., Müller, B.: LBL preprint 12471 (1981)
21. Vasak, D., Stöcker, H., Müller, B., Greiner, W.: Phys. Lett. **93**B, 243 (1981)
22. Stöcker, H.: Pion, Light Fragment and Entropy Production in Nuclear Collisions, LBL preprint 12302 (1981)
23. Gyulassy, M., Kaufmann, S.K.: Phys. Rev. Lett. **40**, 298 (1978)
24. Sandoval, A., Stock, R., Stelzer, H., Renfordt, R., Harris, J., Brannigan, J., Geaga, J., Rosenberg, L., Schroeder, L.S., Wolf, K.L.: Phys. Rev. Lett. **45**, 874 (1980)

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