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Гудима К.К., Тонеев В.Д. **E2 - 12624**

Экстремальные условия, достигаемые в ядерных соударениях при высоких энергиях

На основе каскадной модели обсуждаются величины сжатий и температур, достигаемых при центральных столкновениях ядер, в приложении к возможным фазовым переходам ядерной материи в состояние π-конденсата и в кварковую фазу. Показано, что имеющиеся уже сейчас в Дубне и Беркли пучки тяжелых ионов представляются перспективными для поиска сигналов необычных состояний ядерной материи в лабораторных условиях.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна 1979

Gudima K.K., Toneev V.D. **E2 - 12624**

Extreme Conditions Reachable in High-Energy Heavy-Ion Collisions

Based on the cascade model, we discuss the values of compressions and temperatures attainable in head-on heavy-ion collisions in application to the problem of possible phase transitions of nuclear matter into the π-condensate state and quark phase.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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High energy collisions of two nuclei provide a possibility for the study of hot and highly compressed nuclear matter. Under such conditions, the excitation of pion and quark degrees of freedom seems to be important. In particular, the theory predicts that at not very high temperatures and for nuclear matter compressed two-three times as large as the normal nuclear density, a transition to π-condensate and even the formation of the density isomer are possible [1-4]. At higher compressions and temperatures a transition into a quark phase may occur [5-6]. The search and identification of such phenomena raise many problems. In this work we try to clarify only one of them: Whether it is possible with available beams of accelerated ions to attain extreme conditions under which the phase transitions in question can be realized.

The answer appeals to the study of space-time evolution of the compressed zone and temperature field. In principle, this information may be gained by using the hydrodynamic or kinetic approach. The ratio of characteristic parameters (like the range of NN-interaction, mean free path of a particle in a nucleus, nuclear sizes) which describe the dynamics of nuclear interaction is small but not much smaller than unity as necessary for the conditions of "pure" application of either approach. Therefore, it is impossible a priori to justify a choice of one of them. Moreover, both the approaches give rise to similar results for a number of measured characteristics [7]. The results listed below are found in the framework of the kinetic approach realized by means of the Dubna version of the cascade model of nuclear reactions [8]. The model describes the simultaneous development of intranuclear cascades in both colliding nuclei. It takes into account the processes of production and pion absorption, depletion of the nuclear matter in the course of the intranuclear cascade, the Pauli principle, effects of the relativistic contraction. The model is presented in detail in paper [8].

Figures 1 to 3 demonstrate to what extent the cascade theory agrees with experiment in inclusive distributions of particles for the main reactions discussed in what follows. The model predicts correctly general trends of distributions and the order of magnitude for cross sections. The excess in the calculated number of protons with energy $T \leq 50$ MeV is due to the neglect of the final state interaction of nucleons, leading to the emission of fast complex particles. For the incident beam energy T_0 =0,25 and 0.40 GeV/nucleus, when π-meson production can be neglected, the agreement with experiment is, in general, the same as for other versions of the cascade model. (cf. ref. [7]). The inclusive proton spectra from relativistic (T_0= 2.1 GeV/nucleon) heavy-ion collisions have not been considered in the framework of the cascade model. The comparison of our model with experiment for other characteristics is given in refs. [8,11-13]

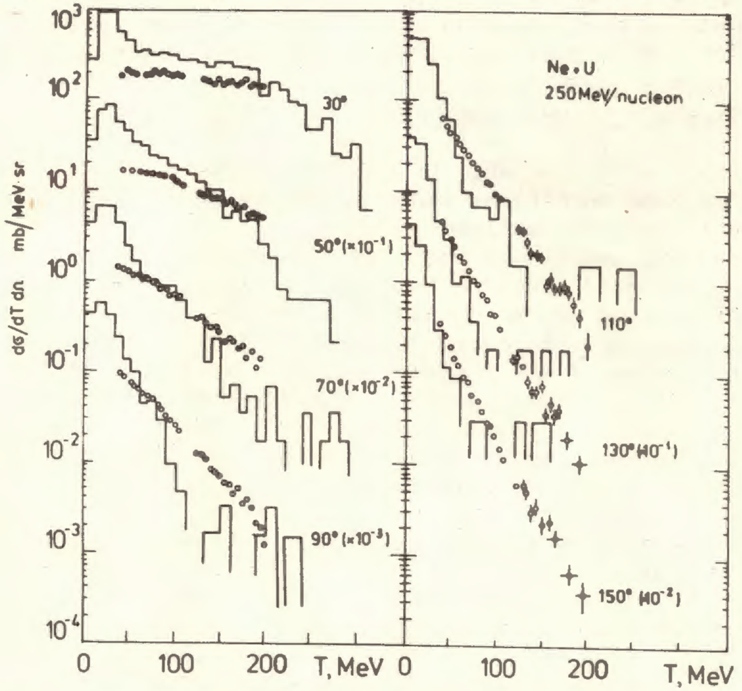

Fig. 1. Inclusive spectra of protons produced in the Ne+U reaction at T_0=0.25 GeV/nucleon. The histograms are calculated within the cascade model. The points are from experiment [9,10].

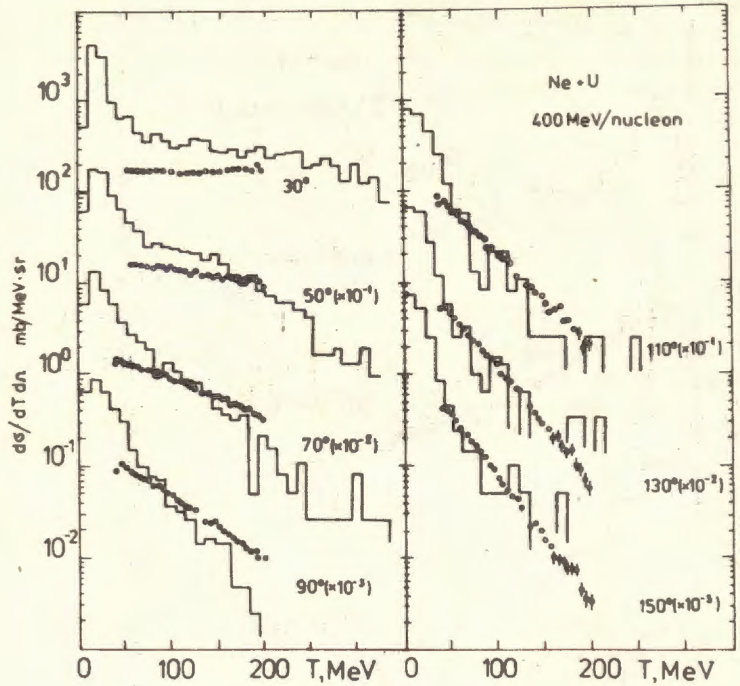

Fig. 2. Inclusive proton spectra for the Ne+U collisions at T_0 = 0.4 GeV/nucleon. Notation is the same as in Fig.1.

Now let us turn to the discussion of the dynamics of nuclear interaction. The cascade model is based on a Boltzmann type equation. Its solution by the Monte-Carlo technique in the lab. reference frame defines the one-particle distribution function $f(\vec{r},\vec{p},t)$. So, the particle number density in the lab. frame is

$$\rho(\vec{r},t) = \int d^3p\, f(\vec{r},\vec{p},t).$$

Because of the specific nature of the Monte-Carlo method, the solution obtained should be averaged over some volume element ΔV around the point $\vec{r}$. We shall consider only the head-on collisions of nuclei. Then the problem possesses the cylindrical symmetry, and the whole range of interaction may be divided into the cells with step Δz along the axis of nuclear collision z, and cylindrical surfaces with

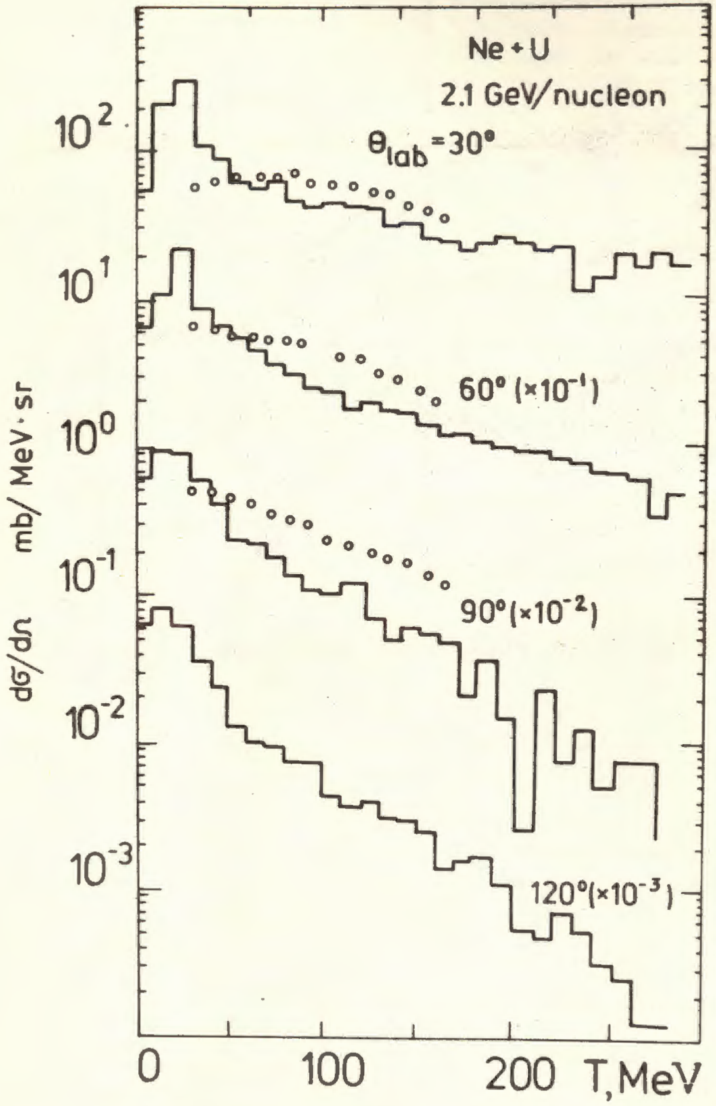

Fig. 3. Inclusive proton spectra for the Ne + U collisions at T_0 =2.1 GeV/nucleon. Notation is the same as in Fig. 1.

step ΔR along the radius. The details of this procedure are described in ref.[14]. The same paper presents the analysis of space-time evolution of the compressed zone (i.e., the behaviour of the function $\rho(\vec r,t)$) formed in inelastic nucleus-nucleus collision.

However, to find a "true" compression ratio, it is necessary to pass to the rest system of every volume element. Therefore, for each s-th cell we defined the current four-vector $j_s = \{\rho_s(\vec r,t), \vec j_s(\vec r,t)\}$, and then go to the system with $\vec j_s - \vec v_s\rho_s = 0$, i.e., $j_s = \{n_s(\vec r,t), 0\}$. The particle-number density in the rest system, $n_s(\vec r,t)$ is related to $\rho_s(\vec r,t)$ by the well known formula

$$n_s(\vec r,t) = \rho_s(\vec r,t)/\gamma_s ,$$

where

$$\gamma_s = 1/\sqrt{1-v_s^2}$$

and

$$\vec v_s = \frac{1}{\rho_s(\vec r,t)} \int_{\{\Delta V_s\}} d^3p\, \vec v f(\vec r,\vec p,t)$$

is the velocity of the volume element ΔV_s (integration runs over the momenta of all particles of the selected cell in a given "sample"). Like the function $\rho(\vec r,t)$ (see ref.[14]) the compression ratio $n(\vec r,t)/n_0$ (here $n_0 = 0.175$ fm^{-3}) with increasing interaction time reaches a maximum, n^{max}, at some point and then rapidly decreases: the expansion stage of interaction. It is interesting to note that for a fixed combination of colliding nuclei the dependence of the maximal compression ratio for nucleons, n_N^{max}/n_0, on kinetic energy T_0 is saturated as early as $T_0 \simeq 1$ GeV. Such a behaviour is a direct consequence of the depletion effect. Experimentally, this manifests itself in the "saturation" of the energy dependence of the mean multiplicity of g-particles (in terms of the photoemulsion works). A similar regime for hadron-nucleus interactions sets in at considerably higher energies $T_0 \simeq 3$-5 GeV[5]. The n_N^{max}/n_0 in the saturation regime grows in magnitude with increasing the projectile size, and for heavy ($A \geq 100$) targets amounts to about 2.5:3.0:3.4:4.0, respectively, for α-particles, oxygen, neon and argon ions.

The cascade model neglects the final-state interaction between cascade particles, i.e., pions and fast nucleons are treated as a mixture of the ideal relativistic gas of Boltzmann particles which is not, in general, at thermodynamic equilibrium. Using the same procedure of the subdivision, one can estimate the four-vector of energy-momentum $\{\mathcal{E}, \vec{\mathcal{P}}\}$ for the cell s , where

$$\mathcal{E}(\vec{r},t) = \frac{1}{\rho_s(\vec{r},t)} \int_{\{\Delta V_s\}} d^3p\, E\, f(\vec{r},\vec{p},t)$$

$$\vec{\mathcal{P}}(\vec{r},t) = \frac{1}{\rho_s(\vec{r},t)} \int_{\{\Delta V_s\}} d^3\vec{p}\, \vec{p}\, f(\vec{r},\vec{p},t).$$

Going into the rest system for given cell, we find the average energy per one particle, $\epsilon = \gamma_s(\mathcal{E} - \vec{v}_s \vec{\mathcal{P}})$. However, in order to relate ϵ with temperature $\tau = \tau(\vec{r},t)$, it is necessary to admit the existence of local statistical equilibrium at instant t. Then, by using the relativistic Boltzmann distribution

$$f(\vec{r},\vec{p},t) = \frac{n \exp(-\sqrt{p^2 + m^2}/\tau(\vec{r},t))}{4\pi m^2 \tau(\vec{r},t)\, K_2(m/\tau(r,t))}$$

one can obtain the temperature in a given volume by solving the equation

$$\epsilon = \frac{n_N m_N}{n_N + n_\pi} \cdot \mathcal{R}\left(\frac{m_N}{\tau}\right) + \frac{n_\pi m_\pi}{n_N + n_\pi} \cdot \mathcal{R}\left(\frac{m_\pi}{\tau}\right),$$

where $\mathcal{R}(x) = 3/x + K_1(x)/K_2(x)$, $K_i(x)$ is the second-order Bessel function, m_N and m_π are, resp., nucleon and pion masses.

To have imagination on the order of magnitude, we show in Fig. 4 the time dependence of temperature for the system Ne + U. The averaging was performed over the whole interaction range, and hence, the temperature corresponds to all the cascade particles regarded as a unique excited system (fireball). It is seen that $\frac{\partial\tau}{\partial t} \simeq 0$ at energies $T_0 \approx 0.25$ and 0.4 GeV/nucleon, and one can speak about an approximate establishment of the statistical equilibrium in the system. In this case, the equilibrium value of τ calculated within the microscopic approach, is close to that one following from the phenomenological (macroscopic) model of the nuclear fireball [16]. For $T_0 \geq 0.8$ GeV/nucleon the system temperature changes essentially during the interaction time and its final value differs greatly from the value calculated within the fireball model. An attainment of the statistical equilibrium in high-energy nuclear collisions is considered in more detail in paper [17]. Here we would like only to note that for all the considered combinations of colliding nuclei the formed system is heated (locally) very quickly and achieves the maximum temperature before the time moment when the maximum compression ratio for nuclear matter is reached.

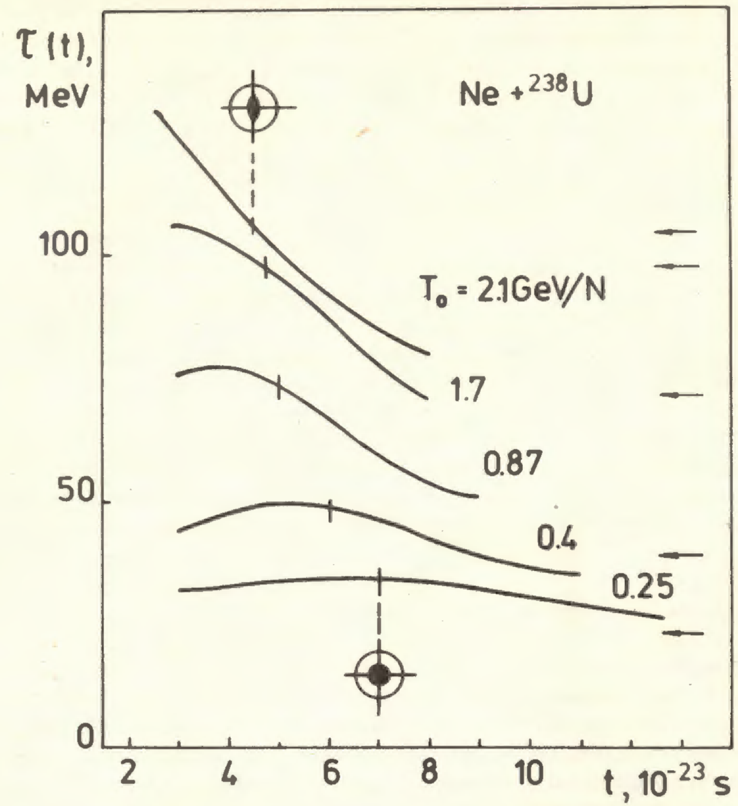

Fig. 4. The time dependence of temperature in the system formed in head-on Ne + U collision at the projectile energy T_0. The values of temperature estimated within the phenomenological fireball model [16] are shown by arrows. The time moment corresponding to the coincedence of the centers of colliding nuclei is shown, as well.

So, the based upon the cascade model calculation scheme presented allows us to trace the evolution of the compressed zone and temperature field when the projectile travels the target-nucleus. Local values of compression and temperature provide an estimate for how the realized conditions are close to the critical point for a phase transition. As an example, we show in Fig. 5 the phase diagram for the transition to the π -condensate state. Among two theoretical curves for the dependence of the critical temperature of the π -condensate on the baryon density n_B, the curve $\tau_c(n_B)$ calculated by Bunatyan [19] seems to be more reasonable. Though it takes into account a lot of extra factors as compared to earlier work [18] the main difference in the final results is due to the fact that in paper [18] the form factor of meson-nucleon scattering is put unity. From the dynamical phase curves $\tau(n_B^{max}, t)$ drawn in the same diagram, it is seen that the increase of the projectile energy from 0.25 to 2 GeV/nucleon does not improve substantially the condition for the π -condensate: Some gain in the compression value is counterbalanced by a noticeable increase in temperature. However, the use of the argon-ion beam available at Berkeley allows us to "look in" the region of the π -condensate restricted by the more strong conditions of Bunatyan [19]*. To form the π -condensate it is important that the time to find out the system in this phase, Δt, would be large enough as compared to the relaxation time, t_R. According to ref. [18] $t_R \approx 10^{-23}\frac{n-n_c}{n}$ s, Galitskii and Mishustin estimated $t_R \approx 1.2 \cdot 10^{-23}$ s if the relative exceeding in density over the critical value n_c is larger than 0.2 [21]. As is seen from Fig.4, for the reaction Ar + Ca (0.5 GeV/nucleon) in the least favourable case, when for $\tau_c(n_B)$ the curve B is taken, the time of the system to be in the π -condensate phase Δt t_R.

Note that the linear dimensions of the cell are chosen to be about 1 fm in accordance with theoretical estimations following from the values of the characteristic momentum of the pion condensate [18,19]. Twice increasing of the linear dimensions of the cell does not influence practically the final results. This estimate does not take into account the role of surface effects, while the π -condensate "germ" is formed, which may, in principle, compensate the energy gained at the transition of the nucleus into the π -condensate phase.

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* It is noteworthy that the dynamical phase curves, calculated within the hydrodynamical approach, are characterized by a qualitatively different behaviour: The system is compressed and heated simultaneously, then the expansion stage occurs during which the system is cooled.

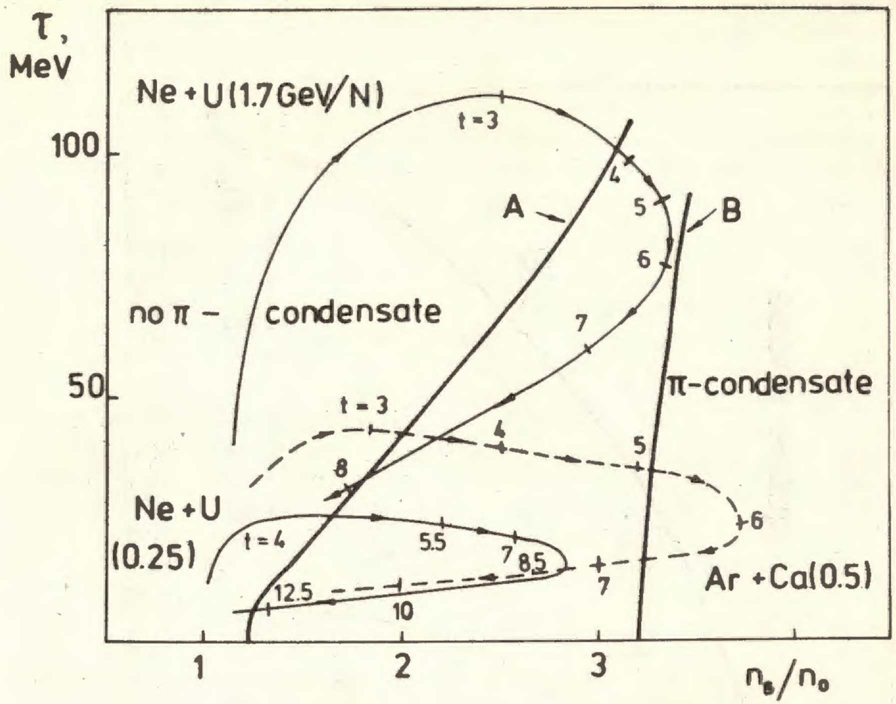

Fig. 5. The phase diagram for a transition into the pion condensate. The curves A and B are critical temperature $\tau_c(n_B)$ from the papers of Ruck-Gyulassi-Greiner [18] and Bunatjan [19] resp. Dynamical phase trajectories are calculated for head-on collisions of nuclei shown in the Figure. The time along trajectories are given in units of 10^{-23} s.

In Fig.6 the phase diagram is drawn for the transition of nuclear matter into the quark phase. One should keep in mind that for this phenomena the temperature plays different role as compared to the transition into the pion condensate. The heating of the system destroys the ordering of the π -condensate. The hadron-quark transition requires a high energy density, and this may be achieved by compressing the system or by its heating at a fixed volume. Therefore the functional dependence $\tau_c(n_B)$ is completely different for these two phenomena.

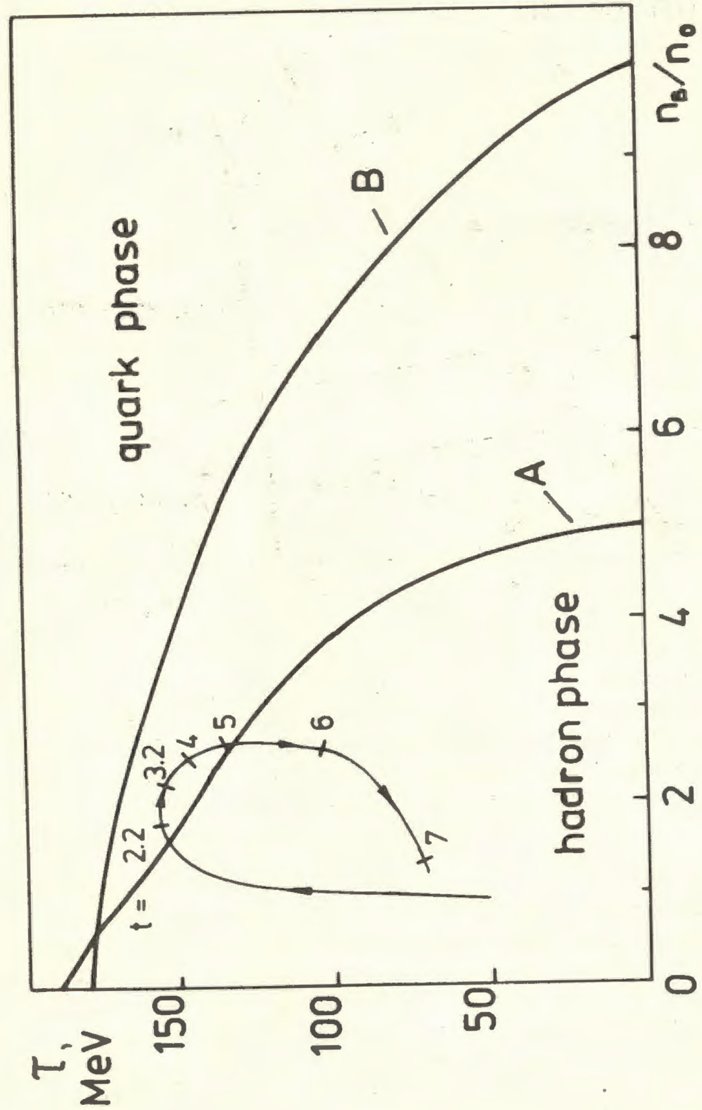

Fig. 6. The phase diagram for hadron-quark phase transition. The critical temperature curves are calculated by Chin[22] (A) and Kapusta[23] (B). The dynamical phase trajectory is shown for head-on ^{16}O + Ag collisions at T_0 = 3.6 GeV/nucleon.

One should not overestimate the accuracy for calculations of the critical temperature of the transition of nuclear matter from the nucleon to quark phase. Theoretical semi-quantitative estimations of $T_c(n_B)$ are highly sensitive to the quark-bag parameters in terms of which the quark state of nuclear matter is described (compare two curves in Fig.6). Besides, it is doubtful whether the perturbation theory can be applied to the problem under consideration (for detais see refs. [22,23]).

However, in spite of all the uncertainties, it is seen that the dynamical phase curve in Fig. 6 for the reaction induced by the oxygen ions with energy 3.6 GeV/nucleon attained at the Dubna accelerator comes rather close to the region of phase transition. The conditions for quark-hadron transition at the Berkeley energy $T_0 \simeq 2$ GeV/nucleon will, naturally, be less favourable than those at the Dubna energy, though the values of the nucleon compression are practically the same. This conclusion differs from the predictions made in [22,23] that this transition is possible even at $T_0 \simeq 1.4$-2. GeV/nucleon. The difference is mainly due to the fact that this work takes into account the finiteness of the interacting system.

Therefore, at present there are possibilities for the search of signs of unusual states of nuclear matter in laboratory conditions. Beams of Ar and Fe ions with energies 300-500 MeV/nucleon available now at Berkeley seem to be more relevant to search the phase transition of nuclear matter in the π-condensate state. High energies of bombarding ions reached at Dubna are hopeful to reveal new phenomena connected with the possible formation of a quark phase of nuclear matter.

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