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THEORY AND PHENOMENOLOGY OF THE QCD VACUUM

7. MACROSCOPIC EXCITATIONS

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НОВОСИБИРСК

A B S T R A C T

This preprint contains discussion of highly excited matter with the energy density ε of the order of 1 GeV/fm^3 at which normal hadronic matter undergoes transition to the so-called quark-gluon plasma phase. In section 7.1 we discuss some questions of the asymptotically dense matter ($\varepsilon \to \infty$), which is considered in perturbative context. In section 7.2 we consider the fate of instanton-type fluctuations in matter, and than in section 7.3 proceed the results obtained in recent lattice calculations. The section 7.4 is devoted to phase transitions, while in section 7.5 we briefly discuss various applications of the macroscopic approach to high energy hadronic collisions, with emphasis on heavy ion experiments.

7. MACROSCOPIC EXCITATIONS OF THE QCD VACUUM

Let us start the introduction to this chapter with the explanation of its title. There are two standard approaches toward investigation of the structure of any kind of matter, based on micro- and macroscopic considerations. The former case deals with microscopic excitations, say the phonons in solids. Usually one (or few) of them are considered in some (arbitrarily large) normalization volume V. The corresponding excitation energy density $\varepsilon \sim O(1/V)$ is small, so one may neglect its backward influence on the matter properties. In chapters 4, 5 we have made such type of analysis, in which infinitesimal probes (external currents) have affected the unperturbed vacuum.

Quite different approach is used in the latter case, in which finite excitation energy density is considered. Moreover, its values of interest are those comparable to the energy density characteristic for matter under consideration. In the problem considered in this paper it is the energy density generated by nonperturbative field fluctuations in the QCD vacuum, considered in section 1.3. Evidently, in such case the matter is essentially rearranged, under certain conditions even qualitatively. In such approach it is natural to consider homogeneous and macroscopically large systems (in order to get rid off the boundary effects) which are also in equilibrium states (in order not to care about the initial conditions).

The methods for corresponding calculations are well developed, and they provide the thermodynamical quantities which are very useful for understanding of some bulk properties of the system considered. Moreover, we know from the history of physics that sufficiently accurate measurements of black body radiation spectrum together with their intensive discussion by theorists have been the starting point of quantum mechanics. We know now that there were plenty of other known effects which in principle also may have triggered its development. However, they were understood much later because free radiation at finite temperature is among the simplest relevant problems.

Being now faced with so complicated problem as that of QCD vacuum structure we also prefer to start with problems, being as simple as possible. Also the properties of matter under extreme conditions are, of course, of permanent interest to physicists. Therefore, discussion of macroscopic QCD excitations were started, see reviews and conference proceedings [7.1—7.5]. It is not possible to consider these topics in details, so in this chapter we only comment on some progress in the field during last 2—3 years.

In section 7.1 we briefly discuss two points connected to perturbative theory of asymptotically dense matter, the so-called plasmon puzzle and the problem of magnetic screening. In section 7.2 we consider instanton-induced effects in the excited matter and show that (in agreement with expectations [7.17]) they are suppressed relative to those in vacuum. Rather impressive progress have been recently made by numerical lattice evaluation of thermodynamical parameters of hot hadronic matter, which we briefly discuss in section 7.3. At least two phase transitions were predicted to take place on general grounds: two main qualitative features of the QCD vacuum such as confinement and spontaneous breakdown of chiral symmetry (SBCS) are expected to disappear at high enough energy density. We discuss these questions in section 7.4. The last section 7.5 deals with perspectives to produce such macroscopic excitations of the QCD vacuum in laboratory by means of high energy collisions of hadrons and, especially, of heavy ions.

7.1. Some problems related to perturbative quark-gluon plasma

It is intuitively clear that in very dense matter (with excitation energy density ε much larger than that of nonperturbative fluctuations ε_{vac}) one should consider mainly the interaction of quarks and gluons at small distances, being relatively small due to asymptotic freedom. Therefore it seems natural that such matter is close to ideal gas made of these fundamental constituents. Perturbative phenomena in such matter (say, the Debye screening of the charge) are very similar to those in ordinary electromagnetic plasma, so it was called the quark-gluon plasma [7.1]. Although main questions relevant to perturbative effects in this phase of matter were solved at late seventies, some of them remain open and in this section we comments on recent publications.

The first point is the so-called plasmon puzzle. In Refs [7.6, 7.7] (in complete analogy to Gell-Mann—Bruckner analysis of QED plasma) some infrared divergent subseries of diagrams was summed up with the finite result, being of the order of g^3. Including this term the thermodynamical potential Ω (we remind that its natural variables are temperature T and chemical potentials for different flavours μ_F) looks as follows

$$\Omega = -\frac{\pi^2}{45}T^4\left[(N^2-1)+\frac{7}{4}NN_f+N\frac{15}{2\pi^2}\sum_F\left(\frac{\mu_F^2}{T^2}+\frac{\mu_F^4}{2\pi^2T^4}\right)\right]+$$
$$+\frac{N^2-1}{144}g^2T^4\left[N+\frac{5}{4}N_f+\frac{9}{2\pi^2}\sum_F\left(\frac{\mu_F^2}{T^2}+\frac{\mu_F^4}{2\pi^2T^4}\right)\right]-$$

$$-\frac{N^2-1}{12\pi}g^3T^4\left[\frac{N}{3}+\frac{N_f}{6}+\frac{1}{6}\sum_F\frac{\mu_F^2}{\pi^2T^2}\right]^{3/2} \qquad (7.1)$$

where the charge g is taken at typical momenta $\max(T, \mu_F)$. Note, that the last term is not analytic in g^2, explaining the divergencies. Note also, that in [7.6] and [7.7] completely different techniques were used (in particular, Coulomb and covariant gauges), but the results are identical (at this point I acknowledge very helpful conversation with J.Kapusta who later have checked my derivation presented in [7.1] and have found an error in the numerical coefficient at the very last step).

The problem with the last term is that it is uncomfortably large even at very large density and/or temperature. Two alternatives are possible here: (i) all further terms have even larger coefficients and the perturbative series are asymptotic and meaningful only for unrealistically small $g(T, \mu_F)$ or (ii) this phenomenon is specific for plasmon effect and it is really there, while the next coefficients are «normal».

In order to understand what is going on in further terms one may either evaluate them explicitely, or make some nonperturbative calculation. Both approaches were recently considered in literature.

In Ref. [7.9] Toimela have evaluated the next term $O(g^4\cdot\log(g))$ (also connected with the plasmon effect)

$$\Delta\Omega = \frac{N^2-1}{32\pi^2}T^4N\left(\frac{N}{3}+\frac{N_f}{6}+\frac{1}{2\pi^2}\sum_F\frac{\mu_F^2}{T^2}\right)g^4\ln\left(\frac{1}{g^2}\right) \qquad (7.2)$$

note that it is also large and has the opposite sign. Similar phenomena in further corrections were also discussed by Kellman and Toimela [7.9]. Also, comparison with Monte-Carlo data (see e.g. [7.10]) discussed in the next section suggests that there is no large correction to thermodynamical quantities above the transition region and approach to ideal gas expressions is here rather fast.

So, everything point toward the first alternative mentioned above, and one may therefore conclude that there is not much sense in further perturbative calculations.

The second problem, that of magnetic screening, is not important quantitatively, but it is potentially connected with some qualitative effects. As in QED plasma, the static gluomagnetic field is not screened in one loop approximation [7.11], in contrast to gluoelectric one. However, in contrast to QED magnetic field, the gluonic one is selfinteracting because of nonabelian effects, and this leads to specific infrared divergences [7.12, 7.13] (see also comprehensive discussion in reviews [7.1, 7.2]). In high orders of the perturbation series there appear powers of the following factor

$$(g^2T/k_{min}) \tag{7.3}$$

where k_{min} is the infrared cut off. The simplest possibility suggested by Polyakov [7.12] is that radius of magnetic screening is of the order of

$$R \sim 1/m_{mag} \sim 1/(g^2T) \tag{7.4}$$

(we remind that electric Debye length is of the order of $1/(g\cdot T)$). However, in recent work by Kajantie and Kapusta [7.15] somehow different possibility was suggested. Using Schwinger-Dyson-type equations for the propagator they have found

$$m^2_{mag} = \frac{3}{8\pi}\left(\frac{g^2N}{3}\right)^3 T^2 \tag{7.5}$$

Although the equations used are not well grounded (corrections to vertex functions are not included), this result is rather interesting by itself. The point (somewhat of academic interest, though) is that (7.4) implies that all terms in thermodynamical potential starting from the 8-th order in coupling constant are of the same magnitude, so that perturbative approach breaks down completely. The possibility (7.5) implies that the powers of parametrically small parameter $g(T)$ is still present in the series, although in smaller power.

Numerical experiments have also been performed [7.16] in order to solve this problem. Existence of finite sreening length was demonstrated, but of course it is not so far possible to distinguish two possibilities mentioned above by their dependence on the coupling. As for the absolute value, (7.5) is reasonably consistent with data.

7.2. Instantons in matter

Qualitative behaviour of instanton-like fluctuations at finite temperature and/or density was discussed in Ref. [7.17] with the conclusion that they are suppressed by the Debye-type screening. During last few years several authors have considered this problem at much more quantitative level. Two limiting cases are usually discussed: (i) large temperature T and zero chemical potentials, and the opposite case (ii) of zero T and large baryonic charge density.

The former case is studied in much details by Pisarsky and Yaffe (see for rather complete presentation and references review [7.2]).

The classical analog of the instanton solution satisfying periodic boundary conditions at finite Euclidean time $\tau_0 = 1/T$, the so-called «caloron», was previously found by Harrington and Shepard [7.18]. It can be written as follows where $\theta \equiv \pi T$

$$A^a_\mu = \bar{\eta}^a_{\mu\nu}\,\Pi(x)[-i\partial_\nu \Pi^{-1}(x)]$$

$$\Pi(t,r) = \left[1 + \frac{\theta \varrho'^2}{r}\sinh(2\theta r)\right]/[\cosh(2\theta r) - \cos(2\theta t)] \tag{7.6}$$

At small temperatures $T\cdot r \ll 1$ it becomes the ordinary instanton, but with shifted radius $\varrho^2 = \varrho'^2/(1+\theta^2\varrho'^2/3)$.

Evaluation of quantum effects (functional determinants in the field (7.6)) made by Pisarsky and Yaffe have lead to the following result:

$$dn(\varrho') = dn(\varrho')_{T=0}\exp\left[-\theta^2\varrho'^2\left(\frac{2N}{3}+\frac{N_f}{3}\right) - A(\theta\varrho')\left(1+\frac{N}{6}-\frac{N_f}{6}\right)\right] \tag{7.7}$$

where the function A is known only numerically:

$$A(\theta\varrho') \equiv \int \frac{d^3x}{(4\pi)^2}\int_{-1/2T}^{1/2T} d\tau(\partial_\mu \Pi/\Pi)^4 - \int \frac{d^4x}{(4\pi)^2}(\partial_\mu \Pi/\Pi)^4_{T=0} =$$

$$= -\ln\left(1+\frac{\theta^2\varrho'^2}{3}\right) + \frac{0.15468}{[1+0.15858(\theta\varrho')^{-3/2}]^8} \tag{7.8}$$

and it is not in fact very important compared to the first term, producing very strong exponential cut off of the instanton density of the type $\exp(-\mathrm{const}\cdot T^2\cdot\varrho^2)$.

Cool quark plasma was studied by a number of authors, but most of these works are devoted to the calculation of the fermionic determinant for the ordinary instanton configuration of the gauge field. It was pointed out in Ref. [7.22] that although this problem is rather interesting by itself it is not directly relevant because it ignores the modification of the instanton solution in matter (e.g. like that in (7.6)). In principle one should start with the calculation of determinant (or effective action) for arbitrary configuration, and only than look for its minimum. This complicated problem is not so far solved.

However, it is expected (see below) that the account for such «feed back» of the matter does not change the results qualitatively, but only modify some numerical coefficients. Therefore, we present available results for the determinants as some methodical example.

Correct result for the instanton determinant in cool quark plasma was first found by Corvalho [7.19]

$$dn(\varrho) = dn(\varrho)_{\mu=0} \exp(-N_f \varrho^2 \mu^2) \tag{7.9}$$

but only in the limit $\varrho\mu \gg 1$ and in very complicated calculation. (It is sufficient to say that the coefficient 1 in the exponent was expressed as some complicated integral, evaluated only numerically.) Than other approach was suggested by Baluni [7.20], but due to some errors the coefficient in the exponent was twice larger. (I acknowledge here his recent letter in which he present some additional details on this point.) Rather straightforward approach based on construction of complete set of states in plasma in the instanton field was adopted by Abrikosov (Jr) [7.21], and after some hesitations he also obtained the result (7.9).

However, rather simple derivation of the general expression for this determinant is possible [7.22] based on the calculation of the scattering amplitude of matter quarks and gluons on the instantons, with further convolution with the thermodynamical weigths. The result looks as follows:

$$dn(\varrho)_{T,\mu} = dn(\varrho)_{T=\mu=0} \exp[-C(T,\mu)\varrho^2]$$

$$C(T,\mu) = \int \frac{d^3p}{(2\pi)^3 2E_p} \left[\frac{8\pi^2 N_f}{\exp[(E_p-\mu)/T]+1} + \frac{8\pi^2 N_f}{\exp[(E_p+\mu)/T]+1} + \right.$$

$$\left. + \frac{16\pi^2 N}{\exp(p/T)-1} \right]; \qquad E_p^2 = p^2 + m^2 \tag{7.10}$$

where three terms correspond to quarks, antiquarks and gluons. This expression agrees with (7.9) as well as with the exponent in (7.7).

Only in the limit of low density quark plasma one can consider the problem outside the scope of the determinant approximation, and its comparison to (7.9) demonstrates that «feed back» effects on instantons in matter are there. The idea of this calculation was suggested by Corvalho [7.19]: it is possible to expand in powers of chemical potential starting from $\mu=0$ case, or the vacuum state. In lowest nontrivial order it gives

$$\left.\frac{\partial^2 \Omega}{\partial \mu^2}\right|_{\mu=0} = \langle \int dx\, dy\, \Pi_{00}(x,y) \rangle \tag{7.11}$$

where Ω is the thermodynamical potential and $\Pi_{\mu\nu}$ is the polarization

operator in the instanton field, previously calculated in [2.12]. (Note that corrections of some numerical coefficient is needed for this result, as well as for Corvalho calculation.) My final result is the following correction factor are as follows:

$$dn(\varrho)_\mu = dn(\varrho)_{\mu=0} \left(1 - \frac{3}{4} \varrho^2 \mu^2 N_f + \ldots \right) \tag{7.12}$$

which should be compared to expansion of (7.9).

I have also addressed the question whether one may interpret the instanton-induced corrections to the thermodynamical parameters of low density matter as a consequence of some quark «effective mass». For cold quark plasma results follow from (7.12) and for hot case from the expansion of (7.8). It turns out that in these two cases the «effective masses» found are not identical (although very close numerically!):

$$M_{eff} = 2\pi\varrho_c (n_+/N)^{1/2} \begin{cases} \sqrt{2} & T \neq 0,\ \mu = 0 \\ \sqrt{3/2} & T = 0,\ \mu \neq 0 \end{cases} \tag{7.13}$$

which demonstrates that they are qualitatively useful but not strictly defined quantities. Note that instantons produce also interaction between quarks, not only their masses.

The last point in this section is connected with application of the results discussed above for the qualitative prediction of matter parameters at which instanton density is strongly (say, by 50%) modified. Using (7.10) we find the corresponding curve on $T-\mu$ plane, the two limiting transition points are as follows

$$T_{trans} \simeq \frac{1}{\sqrt{3}\,\pi\varrho_c} \simeq 120 \text{ MeV}$$

$$\mu_{trans} \simeq \frac{1}{\sqrt{3}\,\varrho_c} \simeq 350 \text{ MeV}, \qquad n_{trans} \sim 1 \text{ fm}^{-3} \tag{7.14}$$

where we have substituted rather small value of critical instanton radius $\varrho_c = 1/3$ fm. In spite of this, one finds rather modest transition parameters (7.14)!

Completing this section we may say that in agreement with expectations all types of macroscopic excitations lead to suppression of instantons. This effect turns out to be very strong (exponential). Somewhat surprisingly, the numerical coefficients are such that even small instantons are affected by rather dilute matter.

7.3. Evaluation of thermodynamical quantities on the lattice

Numerical studies in lattice approximation provide the unique possibility to study transition from rare hadronic gas (at low level of excitation) to quark-gluon plasma (at very high energy densities).

We have already discussed in section 3.2 how nonzero temperature is taking into account by means of periodic boundary conditions in Euclidean time τ_0 such that

$$T=1/\tau_0 \tag{7.15}$$

In lattice calculations at nonzero T one should take the asymmetric lattice, with different spacing in space and time directions, a_σ and a_β. Respectively there are also two coupling constants in the Lagrangian [7.23]:

$$S(U)=\frac{2N}{g_\sigma^2}\left(\frac{a_\beta}{a_\sigma}\right)\sum_{P_\sigma}\left[1-\frac{1}{N}\operatorname{Re}\operatorname{Tr}(UUU^+U^+)\right]+$$
$$+\frac{2N}{g_\beta^2}\left(\frac{a_\sigma}{a_\beta}\right)\sum_{P_\beta}\left[1-\frac{1}{N}\operatorname{Re}\operatorname{Tr}(UUU^+U^+)\right] \tag{7.16}$$

now one can put $a_\sigma=a_\beta=1$ but two parameters g_σ and g_β are necessary. We really have two independent parameters in the problem under consideration, say the lattice volume and temperature. We remind that energy density and pressure can be expressed in terms of statistical sum Z as the following partial derivatives:

$$\varepsilon=-\frac{1}{V}\frac{\partial \ln Z}{\partial(1/T)}\bigg|_{V=\text{const}};\qquad P=T\frac{\partial \ln Z}{\partial V}\bigg|_{T=\text{const}} \tag{7.17}$$

and using the action (7.16) one finds [7.24]:

$$\varepsilon a^4=6N\langle\frac{1}{g^2}(P_\sigma-P_\beta)-\frac{dC_\sigma}{d\xi}(P_\sigma-P_{sym})-\frac{dC_\beta}{d\xi}(P_\beta-P_{sym})\rangle \tag{7.18}$$

$$(\varepsilon-3p)a^4=6N\frac{d(1/g^2)}{d\ln a}\langle P_\sigma+P_\beta-2P_{sym}\rangle \tag{7.19}$$

where $\langle P_\sigma\rangle$ and $\langle P_\beta\rangle$ are the average values of space and time plaquettes, while $\langle P_{sym}\rangle$ refers to symmetrical lattice. Some complication here is the charge renormalization on the asymmetric lattice. Introducing the asymmetry factor $\xi\equiv a_\sigma/a_\beta$ one may write down the following expressions

$$g_{\sigma,\beta}^{-2}=g^{-2}(a_\sigma)+C_{\sigma,\beta}(\xi)+O(g^2) \tag{7.20}$$

where functions $C_{\sigma,\beta}(\xi)$ are evaluated numerically (note that by definition, $C_\sigma(1)=C_\beta(1)=0$). We refer to original works for details and now proceed to discussion of the results.

From the very beginning of such calculations the transition to correct high temperature limit of pure gauge systems (being just the Stephan—Boltsmann law) was demonstrated.

$$\varepsilon\xrightarrow[T\to\infty]{}\varepsilon_{SB}=\frac{\pi^2}{15}(N^2-1)T^4 \tag{7.21}$$

At Fig.1 we show the results of the works [7.24] for SU(2) gauge group, which display this phenomenon clearly. At Fig.2 it is shown the combination $\varepsilon-3p$, absent in ideal gas made of massless constituents (note, that it is the trace of the matter stress tensor, so scale invariance is relevant here). However, including the nonzero effect of nonperturbative effects in vacuum (which are absent in high density matter) one finds some constant contribution

$$(\varepsilon-3p)\xrightarrow[T\to\infty]{}|4\varepsilon_{vac}|=\text{const}(T) \tag{7.22}$$

Note that data shown at Fig.2 well agrees with (7.22) and provide the value of the vacuum energy density! (So far, in SU(2) theory without quarks.)

Generalization of these results to SU(3) group [7.25, 7.26, 7.38] have shown that in this case transition happens to be with strong discontinuity at the critical temperature

$$T_c\simeq(75-80)\ \Lambda_L \tag{7.23}$$

which is the deconfinement transition. With latest Λ_L (see section 2.3) this corresponds to rather high values about 300—400 MeV. Respectively, the latent heat in this transition (see e.g. [7.38]) also becomes uncomfortably large:

$$\Delta\varepsilon\simeq 7T_c^4\simeq 10\ \text{GeV/fm}^3 \tag{7.24}$$

However, very recent publications [7.39—7.41] which more or less include the contribution of virtual quarks demonstrate that their effects definitely lead to strong decrease of transition parameters. As an example, we show at Fig.3 the dependence of ε on T according to [7.41], where 4-th order hopping parameter expansion was used for $K=0.15$ and 0.20. (It is not easy to put these values into physical quark masses, but anyway the latter one is close to critical K, so in this case quarks are rather light.)

Thus, one should wait for more quantitative analysis of virtual quark effects, but the qualitative tendency found obviously points into correct direction. Finally, we hope that their proper account will return us from high values (7.23, 7.24) to more modest transition parameters, about

$$T_{trans} \sim 150 \text{ MeV}, \qquad \varepsilon_{trans} \sim 1 \text{ GeV/fm}^3 \qquad (7.25)$$

suggested by various phenomenological considerations.

7.4. Phase transitions

Existence of phase transitions at high enough excitation level of the QCD vacuum is suggested by general observations that quark-gluon plasma phase is qualitatively different from the ground state, say confinement and quark condensate (manifesting the SBCS) are absent.

The first more elaborated discussion of the so-called deconfinement phase transition was made in the framework of lattice gauge theory by Polyakov [7.28] and Susskind [7.29]. They have explored the general idea that at high temperature gluonic field becomes in some sense classical, so no place is left to confinement. Further general discussion of the physics of deconfinement transition was made in Refs [7.30—7.32]. Among other problems, they concentrate on the symmetry considerations connected with «centre group» transformations (see below) and on relations with some spin systems. However it is very difficult to say at the moment to what extent these considerations are really relevant for understanding of confinement mechanism. So, after some elementary introduction we proceed directly to results of numerical analysis.

As it was already noted in the preceeding section, thermodynamical quantities evaluated numerically in lattice approximation demonstrate some discontinuities, typical for phase transitions. However, following the general arguments by Landau it is desirable to measure specific «order parameters», which are nonzero only in one of the phases, which reveal more clearly the physical nature of the transitions.

For deconfinement transition the following candidate for this role was proposed, the average value of the so-called Polyakov line which is the Wilson contour in the time direction closed by periodicity.

$$L(\vec{x}) = \frac{1}{N} \operatorname{Tr} \operatorname{Pexp}\left[\frac{ig}{2}\int_0^\beta d\tau A_0^a(\vec{x},\tau)t^a\right] \qquad (7.26)$$

Note that the energy $E(\vec{x})$ of two quarks at distance x can be written as follows:

$$\exp\left(-\beta E(x)\right) = \langle L(0)\, L^+(\vec{x})\rangle \qquad (7.27)$$

and in case of linear confinement $E(x) = k\cdot|(\vec{x})|$ one has

$$\lim_{|\vec{x}|\to\infty} \langle L(0)L^+(\vec{x})\rangle \to \langle L(0)\rangle^2 = 0 \qquad (7.28)$$

The opposite case of finite mass renormalization corresponds to the following result:

$$W \equiv \langle L(0)\rangle = \exp\left(-\beta E_{quark}\right) \neq 0 \qquad (7.29)$$

Thus, measurements of W as a function of temperature can reveal the deconfinement transition.

Note that SU(N) gauge field action is invariant over some additional discrete «centre group» transformations, being just multiplication on powers of $\exp(2\pi i/N)$. However, nonzero W implies violation of this symmetry! That is why it was suggested to connect confinement with this symmetry. There exist specific topological objects on the lattice, the vortexes, which may «condense» and in principle produce phases with violation of the «center group» symmetry. However it is not clear what is their fate in continuous limit. It is much more important that «center group» symmetry considerations are not directly generalized to theories with quarks, see e.g. recent discussion in [7.30]. Therefore, their relevance to confinement problem is not clear (as well as what have to be used as deconfinement order parameter).

At Fig.4 and 5 we present SU(2) and SU(3) data for $W = \langle L\rangle$ [7.44], which indeed demonstrate the behaviour expected for the deconfinement transition. Note that in SU(3) case one again observes the finite jump typical for first order transition. However, this strong discontinuity disappears if one includes at least one quark of mass smaller than some critical value (about 800 MeV) [7.39—7.41], and presumably second (or higher) order transition survives in realistic case with light quarks.

Now we proceed to discussion of another qualitative phenomenon, restoration of chiral symmetry. In section 1.5 we have mentioned close analogies between the QCD vacuum and superconductors, Cooper pair condensation and nonzero quark condensate etc. It is intuitively clear that this analogy should go further and at high temperature vacuum should return to «normal» state.

Early discussions of this phenomenon were usually made in the framework of sigma model, but they are rather unrealistic, say the high temperature phase contains massless baryons. Recently this possibility was discussed again in connection to models with composite quarks and leptons, but this interesting point is surely outside the scope of the present review.

Considering deconfinement transition and that connected with the restoration of chiral symmetry in QCD one inevitably ask the following question: are these phase transitions separate ones or the same transition? Definite answer to it will essentially clarify the interplay between these two main qualitative phenomena in QCD.

The following simple consideration provides the partial answer to it (being a part of theoretical folclore, see e.g. latest Kogut review [3.2] and references therein): confinement of quarks is not possible without SBCS. It was already mentioned in section 6.1 that in the bag model one should pay attention to the fact that quark reflected by the bag boundary changes its chirality, which can happen only in chirally asymmetric vacuum. This consideration does not depend on the model and is quite general, so deconfinement may happen either together or before chiral symmetry restoration in the course of increasing level of matter excitation.

In my paper [7.42] the latter possibility was suggested on the basis of some phenomenological observation that SBCS effects are stronger and start to be important at smaller distances (see their discussion in chapters 5 and 6) than those connected with confinement. In short, it was suggested that first hadrons and only than «constituent quarks» are «melted». Pisarsky [7.43] have formulated some phenomenological arguments which show that these two transitions can not in fact be too much separated. However, in terms of excitation energy density this limitation still allows the difference of about one order of magnitude.

Recent numerical experiments devoted to this point where performed by J. Kogut and collaborators [7.44, 7.45, 1.35] and they have clearly demonstrated that separated phase transitions are quite possible. In particular, quark condensate was shown to appear in lattice QED at some critical coupling, in the model without confinement. Most spectacular are results for SU(2) theory with quark «colour isospin» to be larger than 1/2, in particular 1, 3/2, 2: in this case stronger interaction between quarks lead to much larger excitation level needed in order to «melt» the quark condensate. (Note however, that calculations so far ignore virtual quarks!) Results for «normal» quarks with «colour izospin» 1/2 are shown at Fig.4 and it is seen that in this case two order parameters seem to vanish at different temperatures, $T_{chiral}/T_{deconf}=1.3$

or at energy densities differrent by about one order. (However, the authors are careful at this point and comment that in principle it may be some finite size effect.) In the SU(3) case the observed jump in energy density is so strong (about one order by itself) that both transitions are hidden in it, see Fig.5. We have already discussed in preceeding section that this jump is shown to disappear with more appropriate account for virtual quarks, so it is necessary to wait for some time before realistic QCD can be treated by this method.

Now let me make some comments on predictions of the instanton models of QCD vacuum concerning chiral phase transition. Assuming that instantons dominate in the generation of SBCS (see section 2.6) we find rather specific predictions [7.22]. In particular, there is very peculiar dependence on the number of quark flavours.

With only one light quark in the theory the nonzero instanton density directly leads to quark condensate, therefore there is no transition at all: with matter excitation the condensate decreases strongly, but remains nonzero. Quite different behaviour is expected for $N_f>2$ (the realistic case): the selfconsistency equation for the condensate does not allow for solution with arbitrarily small condensate value because the instanton density is proportional to $\langle\bar{\Psi}\Psi\rangle^{N_f}$. As a result, the finite jump from some critical condensate value to zero is inevitable, which implies the first order transition. Note however, that possible influence of Coulomb attraction (at small distances) and confinement effects (at large ones) may in princible somewhat modify the proposed picture.

Other considerations are connected with the instanton suppression in excited matter, discussed in section 7.2. Taking literally these expressions and again assuming that instantons dominate in SBCS one finds rather low transition parameters (7.14), in contrast to very high ones suggested (for pure gauge systems) by lattice calculations for the deconfinement transition. As mentioned before, it is difficult to imagine that inverse order of these transitions takes place. So one may hope that proper account for virtual light quarks in the QCD vacuum will resolve this discrepancy.

Completing this section we may say that the situation with phase transitions is very uncertain at the moment. At least we have realized that they seem to depend in nontrivial way on parameters of the problem (the gauge group, quark masses etc.) so in future this field seems to produce rather rich variety of different situations and, respectively, the interesting physics.

7.5. Macroscopic excitations in high energy collisions

Applications of the methods borrowed from the macroscopic physics to description of high energy hadronic collisions have rather dramatic 30-year-old history. The main ideas were suggested by leading physicists such as Fermi, Pomeranchuck and Landau, and they have really explained the main qualitative features of the phenomenon. However, cosmic ray data available at that time were not sufficient for quantitative analysis. About a decade ago large accelerators have started, providing a lot of data, and few enthusiasts have revived this activity. However, most physicists were very sceptical about its relevance and such models were treated as «heretical» ones (terminology belongs to P.Carruthers [7.71]), which can give only very rough description of data, at best.

It is true, region-type dynamical models provide more detailed information (say, on energy dependence of the total cross section, the shadow effects etc.), while various quark models better account for hadronic structure (say, may also describe fragmentation phenomena). It is also true that excited systems available so far are not large enough in order to make macroscopic arguments quite convincing.

However nonperturbative QCD is very complicated and we are not in the position to confront all this detailed information with the underlying fundamental theory. The situation with macroscopic models is different: the equation of state of «excited vacuum» at given temperature is about the simplest imaginable quantity, and there are methods for its evaluation from first principles. Therefore, experimentalists are now asked not only to go to higher energies, but also to study conditions at which this energy can most effectively be «dissipated» into vacuum degrees of freedom. The perspectives of this new direction of hadronic physics are rather exciting and are widely discussed. The question whether macroscopic excitations are experimentally feasible and some remarks on general goals of this approach are postponed till section 8.7.

In this section we assume that macroscopic excited system is indeed produced in high energy collisions of heavy ions (or even protons) and consider various «effects» most useful for its investigations.

As it became traditional, we start with general discussion of space-time picture of ultrarelativistic collisions on the two-dimentional plot time-longitudinal coordinate x, see Fig.6. As in [7.1, 7.3, 7.4] we emphasize that the energy density (in the CM frame) is very high at the collision point due to Lorentz contraction

$$\varepsilon_{max} \simeq 2\gamma_{CM}^2 \varepsilon_0 \qquad (7.30)$$

where $\varepsilon_0 = m_N n_0 \simeq 0.15$ GeV/fm^3, n_0 being the nuclear density (ε_0 is several times larger inside the proton). However, due to small time of flight and the asymptotic freedom of QCD only some rare hard collisions take place during this stage of the process.

However, as the system expands further it may happen that it goes through the stage of local thermal equilibrium which is characterized by the condition

$$l \ll L \qquad (7.31)$$

where l is mean free path of constituents and L is the system dimensions. The two hypersurfaces at which this condition is violated (see Fig.6) are called «relaxation» and «breakup» ones. Before the former one the system can be considered as a set of quasifree «partons» while after the later one it consists of nearly noninteracting secondaries. This picture is rather different for energies of about 10—100 GeV/N (in Lab frame) where there are chances to produce some «fireball» of excited matter with small relative velocities and «high energies», for which the matter in any case is expanded in longitudinal direction with large velocity gradient (Fig.6 corresponds to this case).

Note that principal condition for existence of thermodynamical region is the existence of two separate time scales:

$$\tau_{relax} \ll \tau_{breakup} \qquad (7.32)$$

Now, why this condition may be fulfilled?

It is not difficult to estimate $\tau_{breakup}$ corresponding to $l \sim L$

$$n \simeq \frac{dn}{dy} \frac{1}{\pi R_\perp^2 \tau_{breakup}} \simeq \frac{1}{\sigma_0 R_\perp} \qquad (7.33)$$

where σ_0 is the cross section for secondary hadrons and $R_\perp$ is the transverse dimension of the system, while dn/dy is the particle number per unit rapidity. As a result one obtains

$$\tau_{breakup} \sim \frac{\sigma_0}{R_\perp} \left(\frac{dn}{dy} \right) \qquad (7.34)$$

Note, that this quantity increases with particle density (and therefore with collision energy) and even for pp collisions at very high energies it may be equal to few fermis.

It is essentially more difficult to estimate τ_{relax} for we do not know the interaction of quarks and gluons at relevant distances well enough. In [7.46] (see also [7.1]) I have used perturbative estimates in lowest nontrivial order for qq, qg and gg scattering and have found gg interaction to be most important which may lead to mean free path in the plasma to be as small as

$$\tau_{relax}\sim 1/[30\alpha_s^2(T)T] \quad (7.35)$$

at temperaturre T. In practice this means that τ_{relax} may be few times smaller than fermi in cases of interest. Note also that in contrast to $\tau_{breackup}$ this quantity decreases with growth of particle density, so the inequality (7.32) may really holds at sufficiently large (pp) collision energy.

Now it becomes clear (see e.g. discussion of the correlators in chapters 4, 5) that the nonperturbative interaction is in this case much more important. However, the statements about the strongest effect in the gluonic case and even the estimate (7.35) suggesting $\tau_{relax}\simeq .2—.3$ fermi seems to hold all right. In this case the existence of two scales (7.32) is just the manifestation of the two scales in the QCD vacuum much discussed above (say, dimensions of «constituent quark» and hadrons) and therefore it may take place even for pp collisions in the few GeV region. Only new experiments will provide reliable data on the rate of mixing phenomena.

I have repeated all these considerations here because the principal importance of two scales (7.32) is often forgotten. Also often raugh estimates of the energy density reached in the collisions are evaluated with some intermediate «formation time» (see e.g. [7.48, 7.49])

$$\tau_{form}\sim 1\,\mathrm{fm}/c \quad (7.36)$$

so that

$$\varepsilon\simeq\frac{\langle E_\perp\rangle}{\pi R_\perp^2\tau_{form}}\frac{dn}{dy} \quad (7.37)$$

In chapter 6 we have already discussed the value of the formation length extracted from $N—A$ collisions and have concluded that it corresponds to «formation time» of constituent quarks, both according to the model used and to particular numerical values obtained for it from the fit. Thus, it is several times smaller than (7.36), respectively ε in (7.37) is essentially underestimated.

Now we turn to completely different type of problems, connected with the question of most convenient experimental signatures of all these phenomena. In his recent talk at Brookhaven 1983 conference [7.3] M. Gyulassy have suggested very nice terminology. All observables are classified as «thermometers», «barometers» and «seismometers», if they are most sensitive to temperature, pressure or fluctuations, respectively. In addition to this list of useful devices let me also add the interferometric «microscope», directly measuring space-time properties of the system producing observable secondaries.

Since «barometers» based on explosion-type expansion are more spectacular, we start with this point. Than we consider «thermometers» which are mainly based on nonequilibrium phenomena.

We remind that applicability condition of hydrodynamics is the same as for thermodynamics (7.31), but in order to apply it one should also fix boundary and initial conditions, as well as the equation of state. The most uncertain point here is of course initial conditions, depending on poorly understood relaxation rate. In [7.46, 7.1] we have discussed two extremes, instantaneous mixing of Fermi—Landau (Lorentz contracted discs) and «scale invariant» initial state defined on the hyperbola

$$t^2-x_\parallel^2=\tau_{relax}^2 \quad (7.38)$$

The latter condition leads to very peculiar expansion: all volume elements are expanded linearly with time and move in straight lines from the collision point:

$$V(x,\ t)=t/x=\mathrm{const} \quad (7.39)$$

In [7.1] interesting resemblance of this regime to Universe expansion was noticed, with (7.38) resembling the Habble's law. For each volume element picture is essentially the same, and pressure from the left neighbour is equal to that from the right one, respectively there is no acceleration. This scaling regime was first considered in Refs [7.76], and recently rediscovered in [7.48, 7.49].

Note, that in reality the total energy is limited, as well as the hyperbola (7.38), and so some amound of energy should be transmitted to fragmentation region where the ideal pressure balance is lost. Another way to see this phenomenon is to consider the law of energy density dependence on time [7.84]

$$\varepsilon(t)\sim t^{-(1+c^2)}$$

(here the equation of state is $c^2\equiv dp/d\varepsilon=\mathrm{const}$). If the energy be fixed, it would correspond just to $\varepsilon(t)\sim 1/t$.

These two types of initial conditions result in different spectra of longitudinal rapidities. In Landau case it is the famous Gaussian:

$$\frac{dn}{dy}\sim\frac{S^{(1-c^2)/2(1+c^2)}}{\sqrt{L(S)}}\exp(-y^2/2L(S))$$

$$L \simeq \frac{4c^2}{3(1-c^4)} \ln \frac{S}{4m^2} \qquad (7.40)$$

(Landau result is slightly generalised [7.69] for $c^2 \neq 1/3$.) In scaling case one evidently has the familiar «plateau» in rapidity distribution. As it was argued above, limited relaxation rate should produce Transition from the former «stopping regime» at relatively low collision energy to «scaling» one at higher energies. As it was estimated in [7.46, 7.1] this transition was expected at approximately ISR energy region, $s=10^2-10^3$ GeV^2. Now SPS collider data are available, and one can really see that in contrast to ISR data (reasonably well corresponding to (7.39), see e.g. [7.75]) they are much more of the «plateau» type.

Although such qualitative agreement with expectations is found, I am rather sceptical about attempts to obtain more quantitative information along this line. The observed rapidity spectrum is not much sensitive to initial conditions. The readers can find more details in recent Refs [7.82, 7.83] devoted to numerical hydrodynamical calculations.

Much more reliable «barometer» seems to be transverse momentum distribution, or even average $\langle p_\perp \rangle$. In principle the observed spectrum is generated by two components, distributions over «thermal» and «collective» velocities. The general method to separate them is rather simple: one should compare spectra for particles of different mass, say pions and antiprotons. Obviously, thermal velocities for heavier particles are smaller. This idea was, for example, used successfully at low energy heavy ion collisions [7.80] with rather convincing arguments for the existence of collective effects.

Zhirov and myself [7.74] have made similar analysis of ISR *pp* data and the conclusion was rather disappointing: thermal distribution describes data «too well», so that no trace of collective flow velocity (at the level $v<0.2$) was seen. The reason for it is not so far clear. It can be that due to «vacuum pressure» effects the driving force is in this case smaller than expected with simplified equations of state [7.74]. In other terms, most of the energy may be used not for kinetic energy, but for production of the large «bag». Another possible explanation [7.60] is that the system is cooled too rapidly from the surface. However, estimates show that hadron «evaporation» rate from quark-gluon plasma is suppressed by some combinatorical factors and this mechanism is not very effective.

However, recent experiments with essentially larger systems and higher energies have produced some results which allow to look at the

problem with greater optimism. Studies at SPS collider [7.87, 7.88] and first results from high energy nuclear collisions obtained by JACEE collaboration [7.89] have shown that the average transverse momentum increases with the particle density dn/dy, see Fig.7. (It would be also desirable to plot local density here, not the global one, but it is not so far done.) Unfortunately, little amount of data for heavy particles is reported, but UA5 group have observed very spectacular increase of the average transverse momentum for kaons [7.87]. (The explanation proposed in this paper relates this fact to charm decay, but it seems to be rather unprobable: the number of charmed quarks should be comparable to that of strange ones!)

Assuming that JACEE effect is indeed due to pressure, we are in rather puzzling situation: hydrodynamics predicts that the larger is the system, the smaller is the pressure gradient and therefore collective velocity, while data suggest another trand! This naturally suggest that the system is not purely hydrodynamical, at least for small systems. As for JACEE effect in absolute magnitude, it reasonably well corresponds to what can be expected (see e.g. report by G.Baim et al. in [7.83]). So, after all our «barometer» indicates something, but it will need some time to calibrate it.

Now we proceed to nonequilibrium phenomena which give us more direct information on earlier stages of the collisions. The general introduction as well as some examples were already given in [7.1], so we do not consider them in details. During last few years many papers were devoted to the subject, but we do not consider them in details here. The problem with them is that they use the same ideas as earlier works (say, [7.46]), in particular ideal quark-gluon gas and perturbative (lowest order) estimates for the production rate of dileptons, new flavours, surface «evaporation» etc. This is good enough for qualitative estimates, but in order to make the predictions more reliable one needs either make some nonperturbative calculations (say use Kubo-like formulae for dissipative parameters in lattice calculations) or phenomenologically analyse some data. Therefore, instead of presenting of a number of curves I prefer to outline a set of qualitative ideas which may turn useful later for the analysis of (future) heavy ion data.

The most helpful phenomenon of this kind is the production of penetrating particles—photons and leptons [7.50], easily going out of hadronic cluster (like neutrinos from the Sun). Note that perturbative evaluation of dilepton production rate [7.46] is also more reliable here, for it is the electromagnetic process. We also comment that quark-hadron duality ensures that global production rate is very insensitive to the dynamical details.

It is not so for the shape of the spectrum, and this point was emphasized by Domokos and Goldman [7.54] and later by Pisarsky [7.43]. Suppose we start with dilute pion gas and consider pion annihilation into e^+e^-. From the inverse process (well studied at colliding beams) we know its amplitude (in vacuum), it is affected by pion-pion interaction and is strongly enhanced at rho-meson mass with some suppression around it due to duality. In the quark-gluon phase there is no confinement and therefore no rho-peak: the spectrum should be continuous. (It looks as if one observes $e^+e^- \Rightarrow \bar{q}q$ directly.) Interesting, that data even at rather low energy (e.g. see [7.54]) clearly demonstrate significant continuous component in the dilepton mass spectrum. It was so striking that such pairs are known as the «anomalous» ones. Thus, looking at relative contribution of resonance and nonresonance dileptons we may have some information on matter properties.

Another qualitative phenomenon connected with the shape of the spectra was found by Zhirov [7.59]. At Fig.8 we show the local slopes

$$T_{eff}^{-1}(E_t) = -\frac{d}{dE_t} \ln \left(E \frac{d\sigma}{d^3p} \right) \qquad (7.41)$$

where the transverse energy is $(p_t^2+m^2)^{1/2}$. Note that data of quite different experiments and reactions point toward existence of some «plateau» at E_t=3—4 GeV. The proposed interpretation [7.59] is that plateau height is just the «initial temperature» of the system while particles in this E_t range are «evaporated» at the mixing stage. (See earlier discussion of such processes in [7.55—7.60].) Energy dependence of this plateau height and the preexponent are shown at Fig.9, they are in reasonable agreement with behaviour expected for simplest «stopping regime». Evidently, existence of this phenomenon at SPS collider energies and its parameters is of interest.

If this interpretation will be confirmed by later studies, it suggests remarkably simple «thermometer» for measurements of the highest (as well as intermediate) temperatures of the process just by evaluation of the transverse momentum slopes at some well defined kinematical regions. Note that for systems containing thousends of secondaries it can be done even for individual events!

Now we proceed to discussion of the contents of secondaries, which is also a «thermometer», but much more complicated. It is probably useful to recollect that original Pomeranchuck idea is that at final stage we have nearly ideal hadronic gas at low temperature $T_{breakup}$, estimated from the slope of spectra at low transverse momenta. This idea reasonably well reproduces the global number of kaons and antiprotons, observed in hadronic reactions. However, looking at this long-standing prediction more quantitatively on the basis of modern data one comes across new problems.

First of all, pions, kaons and antiprotons seen in detectors are in fact the decay products of multiple resonances. Their widths are only about 100 MeV or less, so they mostly decay after $\tau_{breakup}$, at very low density, thus in well known way. Evidently one should exclude from the consideration this trivial stage of the process!

Fortunately, transverse energy distribution for resonances is remarkably identical to that of stable particles, so we enjoy the so-called «thermal» equilibrium. But what about «chemical» one, related to absolute resonanse cross sections?

Unfortunately, it is very difficult to answer this question at the moment. Relevant temperatures are around phase transitions and we do not even understand how to pick up necessary degrees of freedom in order to understand «how many rho-mesons» are there. At smaller T we have hadronic gas picture and Beth-Ulenbeck method (much advocated by Hagedorn) suggests simple picture of ideal gas made of stable particles as well as of resonances. But even this picture contains many uncertainties, say too wide resonances must be somehow cut off. Even in the pion gas one finds some surprises [7.1]: contribution to thermodynamics of «resonant» $I=0$, 1 channels in pion-pion interaction is canceled by the nonresonant $I=2$ interaction (in first order of Weinberg effective Lagrangian)!

At larger T we may try to form resonances out of plasma quarks. In this picture it is natural that, say, direct ϱ/π ratio is about 3, according to the number of spin states. This is indeed what is seen in data (see more in Ref. [6.64]), and these observations may imply that «chemical equilibrium» is indeed absent at breakup stage and contents is governed by larger temperatures than $T_{breakup}$.

Among surprises from the experimental side I should mention here large baryon excess in processes related to gluonic intermediate stage: Ψ and Υ decays and gluon jets. Especially interesting is Υ decay, producing in average nearly isotropic cluster of a dousen of pions, with baryon fraction twice larger than observed under similar multiplicity in pp and e^+e^- cases.

It is not possible to understand all these puzzles without consideration of essentially nonequilibrium physics. It is probably useful to recollect here some widely known examples of the kind. The first one which comes to my mind is the steel tempering. At large temperature a lot of dislocations appear which are frosen by rapid cooling. Recently many similar phenomena are discussed in cosmology: relict monopoles, baryon charge excess (very tiny effect, but leading to our existance) etc.

The key parameter in all these examples is cooling or expansion rate, which has to be compared to typical relaxation time. In hadronic reactions a variety of cases is possible, depending on initial conditions. For example, it is quite obvious that spherical expansion from (point-like) upsilon decay can be more rapid that transverse (cilindrical) one in pp or e^+e^- cases. The energy spectra in this decay indeed have larger slopes, indicating larger decay temperatures or stronger collective effects. Thus, there are arguments that this system breakes up at more dense stage, explaining the baryonic excess. Selecting larger multiplicity we have larger and more «cool» systems, so we predict that here baryon/meson ratio should gradualy decrease. It is interesting to test this prediction experimentally.

Among nonequilibrium phenomena the most interesting ones are those connected with phase transitions. This point was recently adressed by L. van Hove [7.85], see also recent work [7.86]. It was pointed out that under certain conditions (first order transition with sufficiently large latent heat) at the expansion stage some discontinuities may develop, the so-called deflagration and detonation fronts, at which the internal energy of matter is directly transformed to that of collective flow. Deflagration has more chances to take place, and it can be observed by the fixed value of outword collective velocity v_{coll} of matter, which is determined entirely by the equation of state. However, in problems of similar type a lot of instabilities are known to be developed and instead of some cilindrical front one may in fact find a mess of «bubbles» [7.86] and other «supercooling» phenomena.

Finally, let us consider one more interesting «thermometer» (being somewhat similar to ordinary medical one) the production of new flavours. As dislocations in tempered steel, they are produced at the hot stage and then remains there because the rate of their annihilation is negligible. As some experimental input we may consider data for strange quarks (analyzed in recent paper [7.66] including resonance decays) in terms of «suppression factor» $s/u=\lambda_s$ being direct s quarks yield relative to that of u, d ones:

$$\lambda_s = \left(\frac{s}{u}\right)_{direct} \begin{cases} 0.15 & \sqrt{S}=5\ \text{GeV} \\ 0.39\pm.07 & \sqrt{S}=540\ \text{GeV} \end{cases} \qquad (7.42)$$

Theoretical estimates [7.63—7.65] have suggested that in hadronic gas and quark plasma this parameter should be at least one order different. Indeed, ideal gas formulae

$$\lambda_s^{gas} = F(m_K/T)/F(m_\pi/T)\,, \qquad \lambda_s^{plasma} = F(m_s/T)/F(0)$$

$$F(x) = x^3 \int dy\, y^2 [\exp(x\sqrt{1+y^2})-1]^{-1} \qquad (7.43)$$

produce at $T=140$ MeV (gas) $\lambda_s=.07$ while in plasma at $T=200$ MeV one really finds the equilibrium value about 0.4. Estimates presented in the works mentioned show that in plasma phase production of strangeness is sufficiently fast process, so that this number can indeed be reached.

Production of charmed quarks was considered it Refs [7.62]. In this case no equilibrium can be reached, of course, but the cross sections are not too low: say SPS experiments with heavy ions may have about one pair of charmed quarks per event! Note that production rate is more reliably calculated in this case, although experiments are of course much more difficult.

Now we come to our last device, the interferometric «microscope» based on correlations of two identical pions [7.61]. This idea was suggested first by Kopylov and Podgoretsky, and in my work it was shown that results of observations may directly produce Fourier transform of the space-time structure of the pion source. In spite of so important information which can be obtained by this method it was not much used so far, for the probability for two pions to be in usuful kinematical region is rather low. Among the physical results let me mention that it was clearly demonstrated that the source dimensions really increase strongly with multiplicity. However, with future heavy ion experiments with $\langle n_\pi \rangle = O(1000)$ it will become much more practical because the number of $\pi^\pm\pi^\pm$ pairs becomes $O(1{,}000{,}000)$! With this method one may directly control the energy density, (may be even in individual events), make some detailed space and time «slices» selecting pions with particular kinematical region, etc.

Summary of all such ideas can be made as follows. Measuring temperature by «thermometers» and volume by the «microscope» one can obtain information on the equation of state, being the energy density dependence on temperature. Additional test is provided by pressure measurements by «barometers». Of course, this schematic picture will become much more developed with real experiments going on.
measurements by «barometers». Of course, this schematic picture will become much more developed with real experiments going on.

In conclusion, we have rapidly growing field of investigations, motivated by applications of macroscopic methods to high energy collisions. So far, we have only qualitative ideas and some preliminary data, but the perspectives look very promising.

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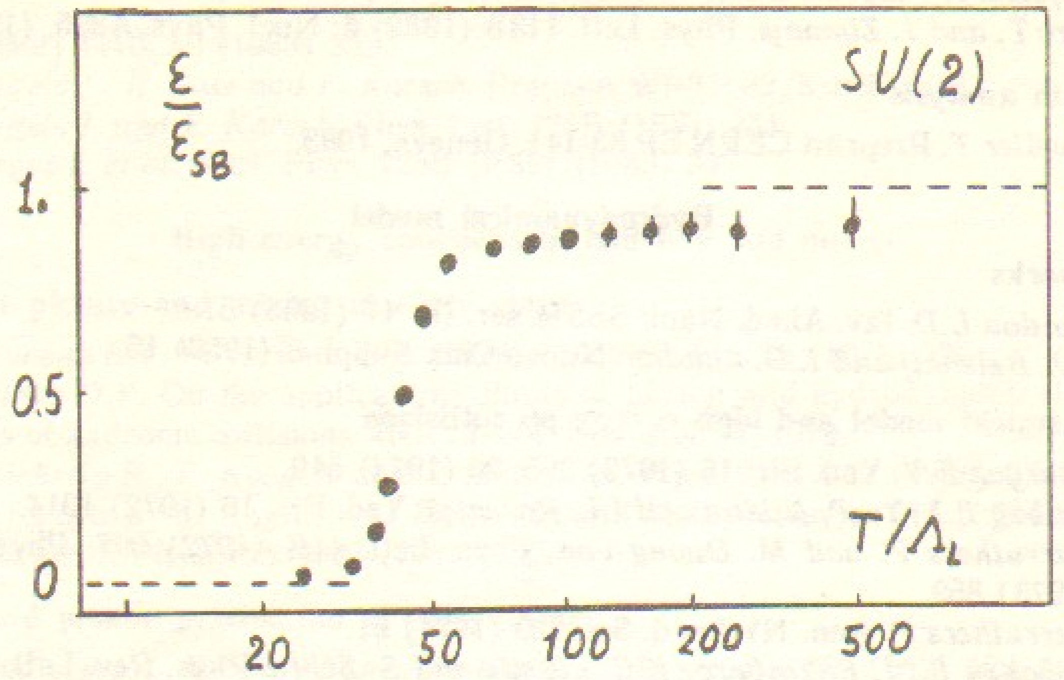

Fig.1. Ratio of the energy density to its asymptotic Stephan Boltsmann value $\varepsilon/\varepsilon_{SB}$ versus temperature T for SU(2) pure gauge theory according to Ref. [7.24].

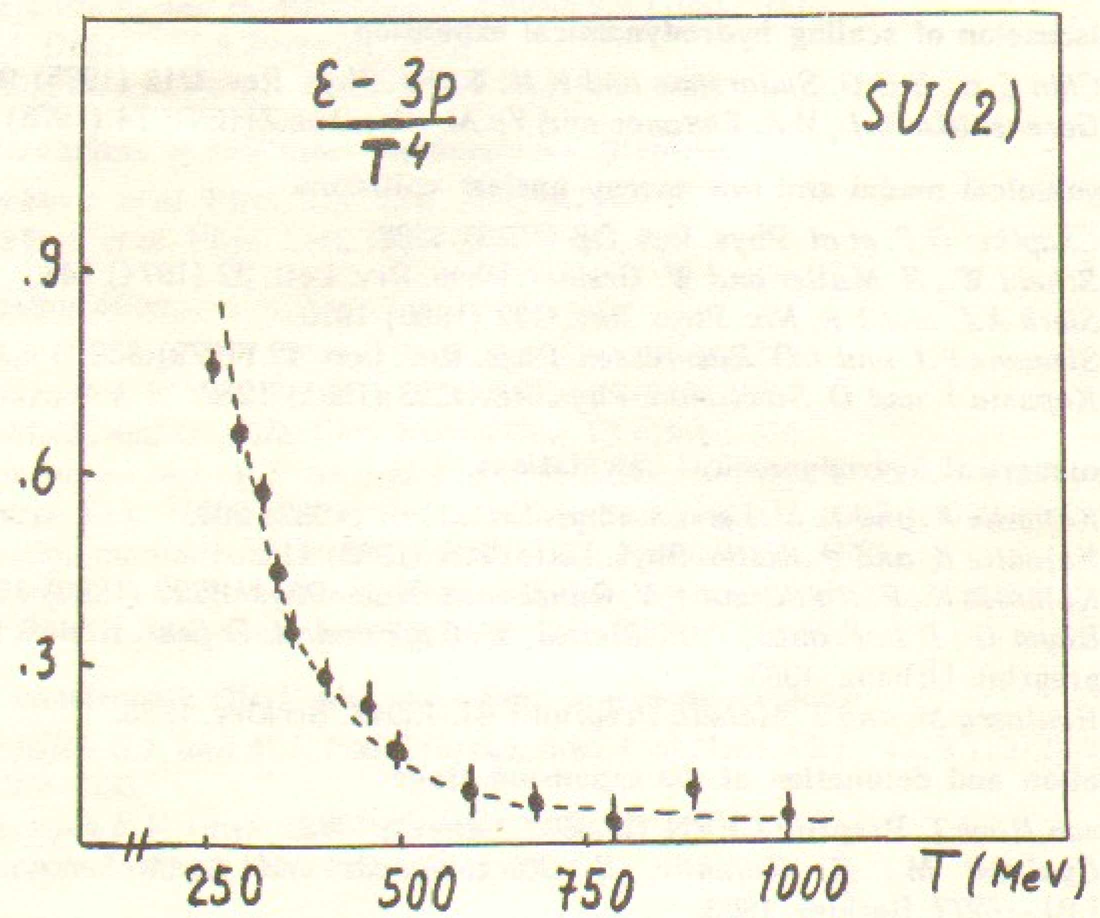

Fig.2. The same calculations as at Fig.1, but with energy density substituted by the combination $\varepsilon - p$ (p is pressure) which vanishes in the ideal gas made of massless constituents. As data shows, this quantity really decreases with T. The dashed line represents some fit of the type const/T^4, which shows that data are consistent with the interpretation of nonzero effect as being due entirely to finite vacuum energy density (see text).

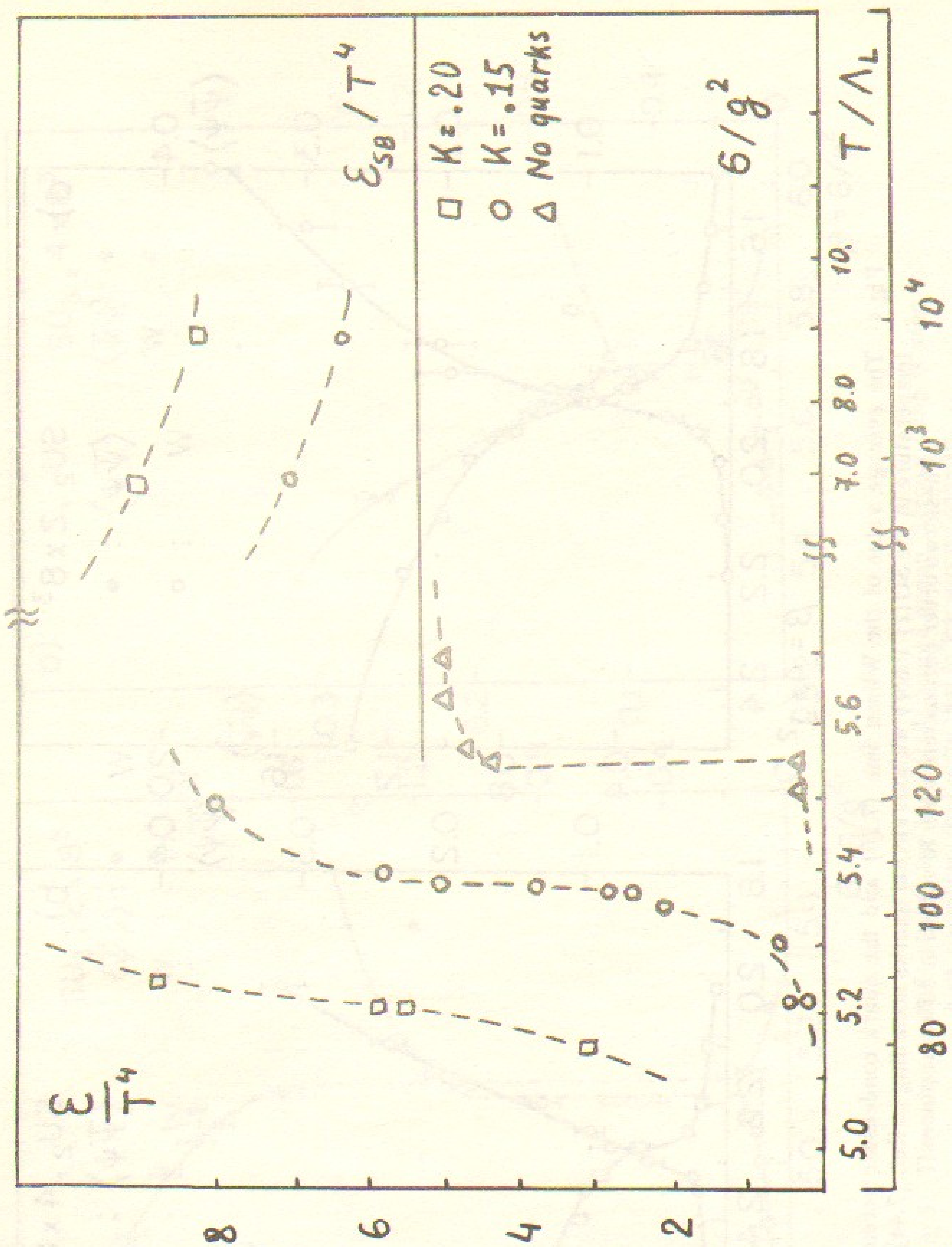

Fig.3. The ratio ε/T^4 as a function of the coupling constant g or the temperature T (in lattice unites of Λ_L). The points correspond to Ref. [7.41] made for the SU(3) gauge group in the fourth order in hopping parameter K at K=.15 and .20, as well as for K=0 (no quarks, note that in this case the temperature scale is not valid). The dashed lines only shown to guide the eyes. It is seen that for lighter quarks the transition temperature decreases, as well as the energy density in absolute unites. Also the finite jump seems to disappear in the presence of sufficiently heavy quarks.

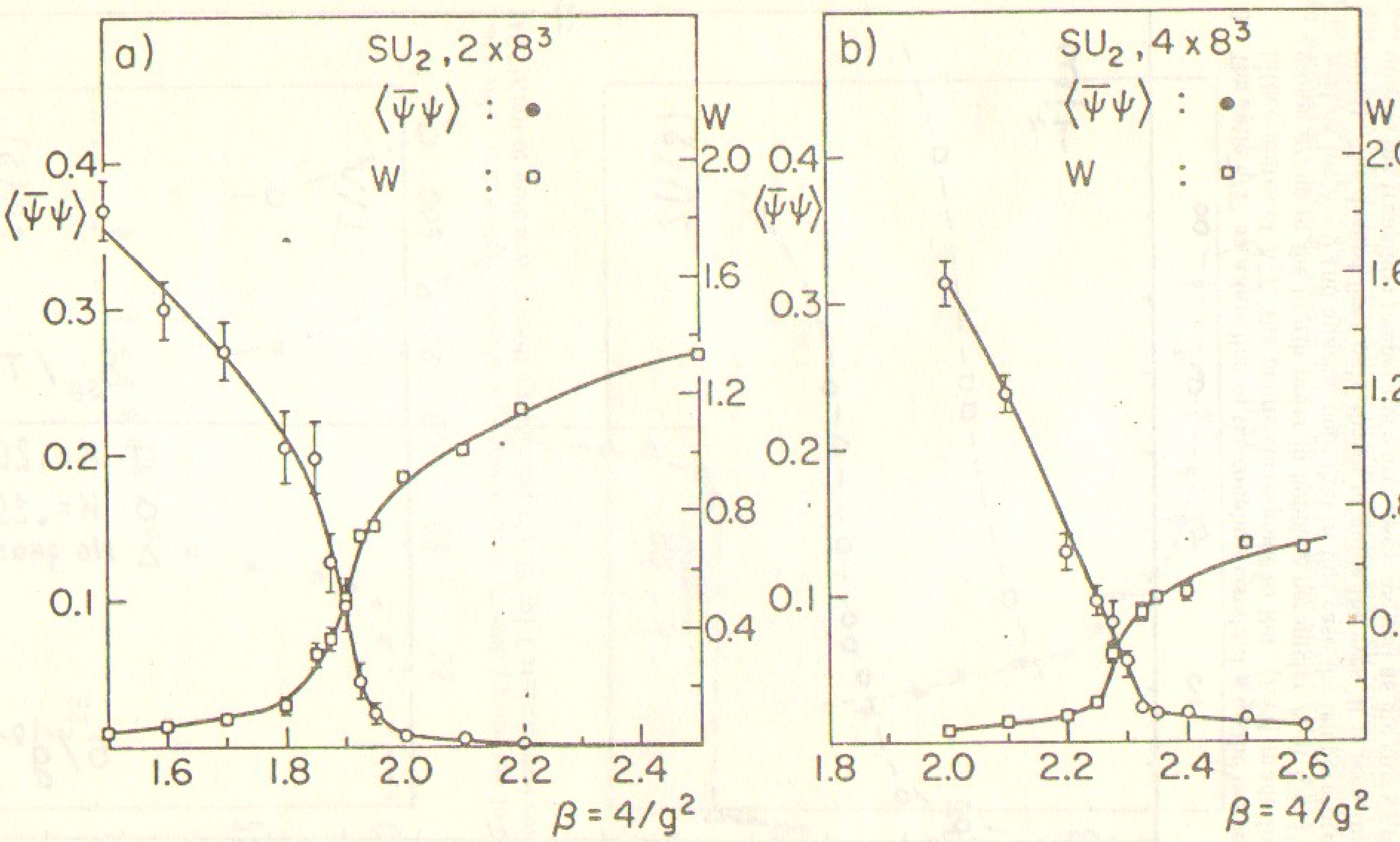

Fig.4. The average value of the Wilson line (7.26) and the quark condensate versus temperature in the SU(2) theory without virtual quarks according to Ref. [7.44]. These two order parameters seem to vanish in different places.

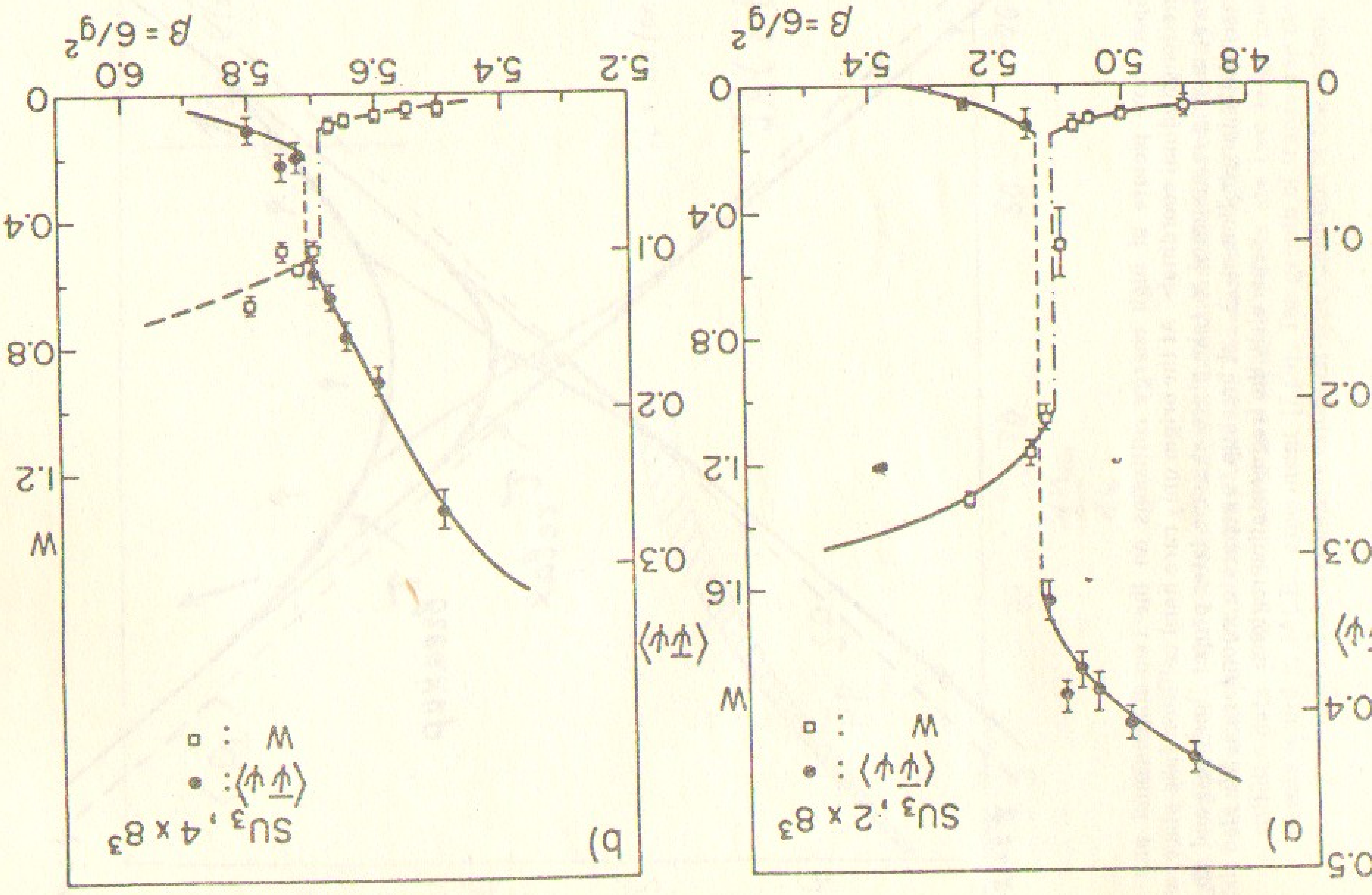

Fig.5. The same as at Fig.4, but for the SU(3) gauge group. Strong jump is seen in both variables, similar to that in the energy density.

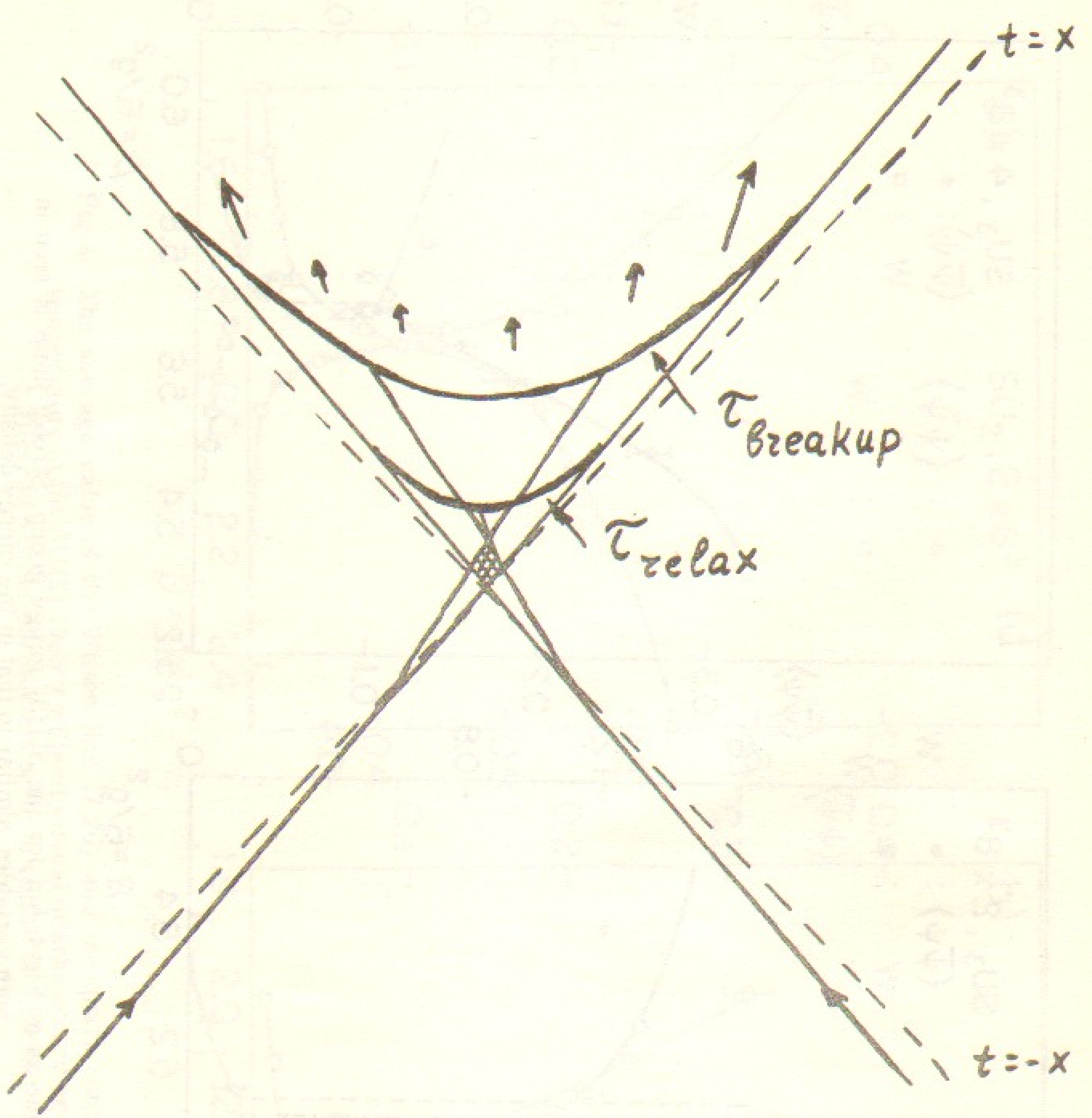

Fig.6. Space-time picture of high energy collisions on the two-dimensional plot time-longitudinal coordinate. At the origin only rare hard collision take place, and local mixing is assumed to take place at some later proper time τ_{relax} and up to essentially larger time $\tau_{breakup}$ of breakup to separate secondaries. We also show separately the fragmentation regions.

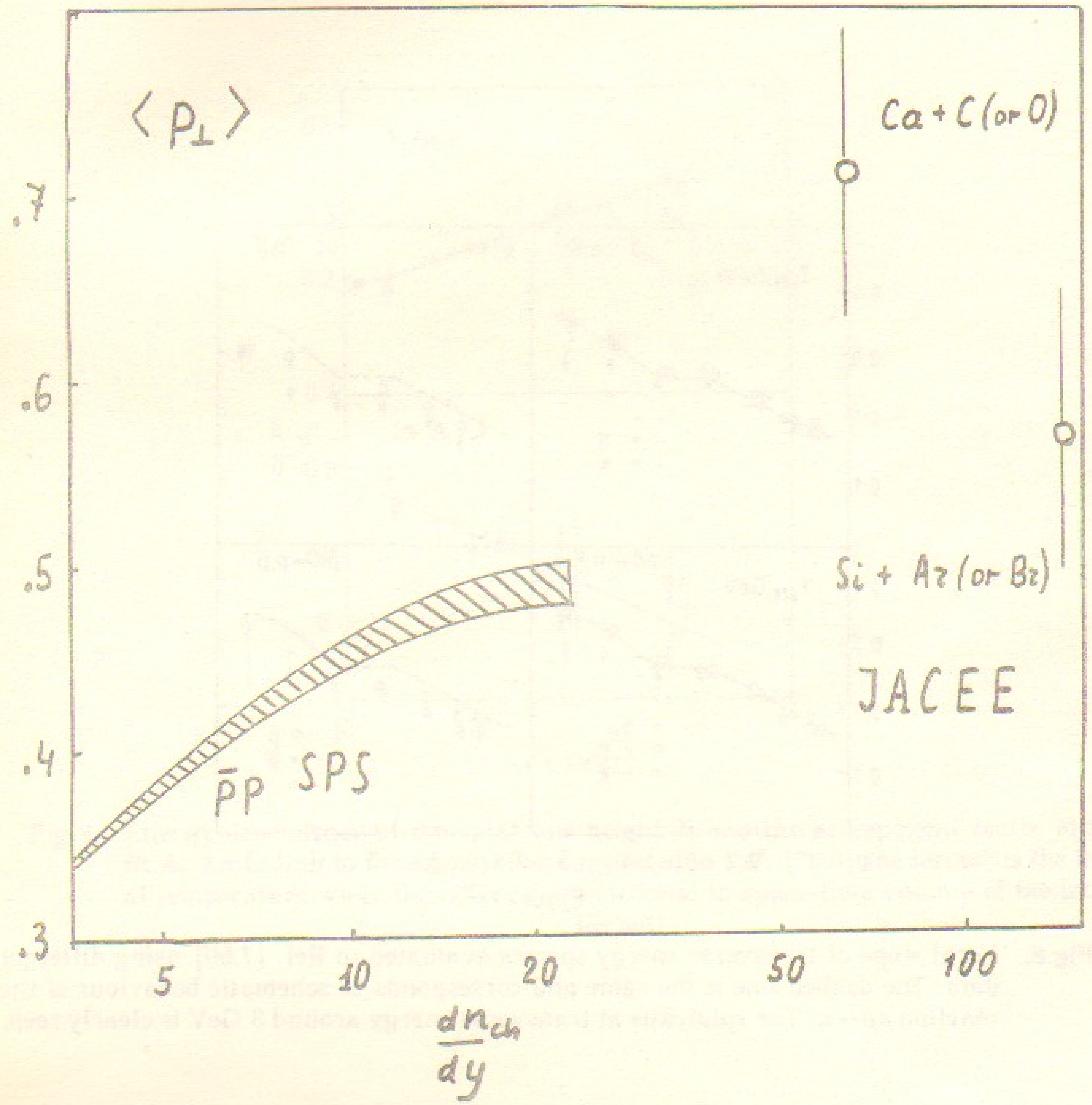

Fig.7. Dependence of the average transverse momentum of pions on density of charge secondaries per unite rapidity interval. The dashed region corresponds to SPS collider data [7.87, 7.88] while two separate points are two large multiplicity events found in JACEE experiment [7.89]. The grouth of transverse momentum may be due to collective explosion-type expansion of the system.

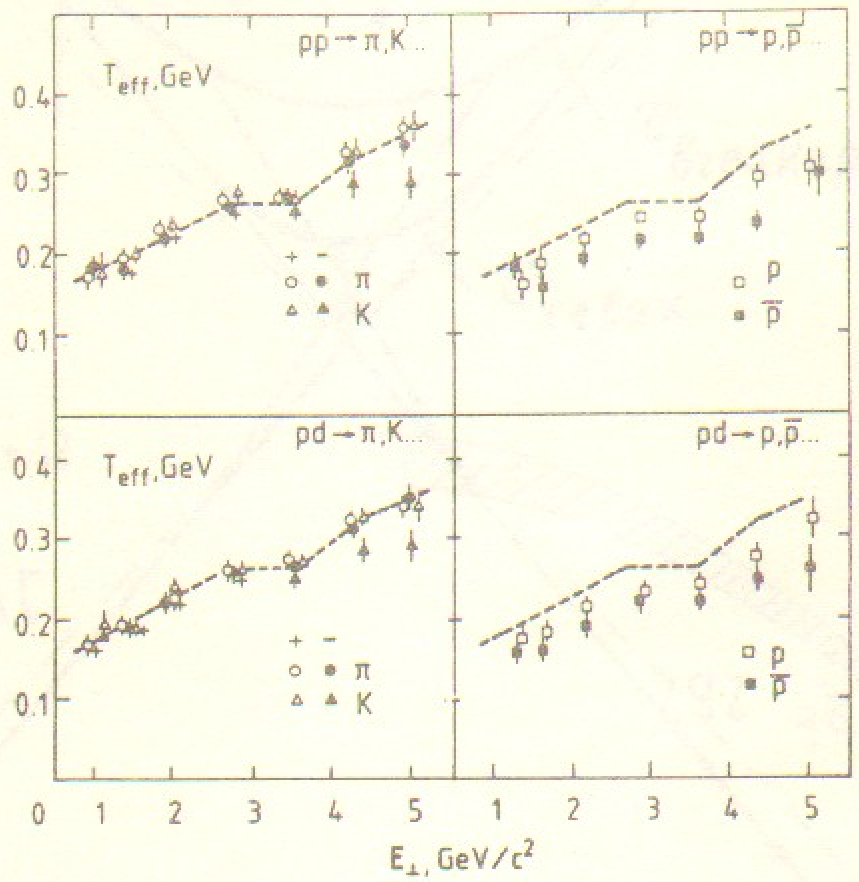

Fig.8. Local slope of transverse energy spectra evaluated in Ref. [7.59] using different data. The dashed line is the same and corresponds to schematic behaviour of the reaction $pp \to \pi$. The «plateau» at transverse energy around 3 GeV is clearly seen.

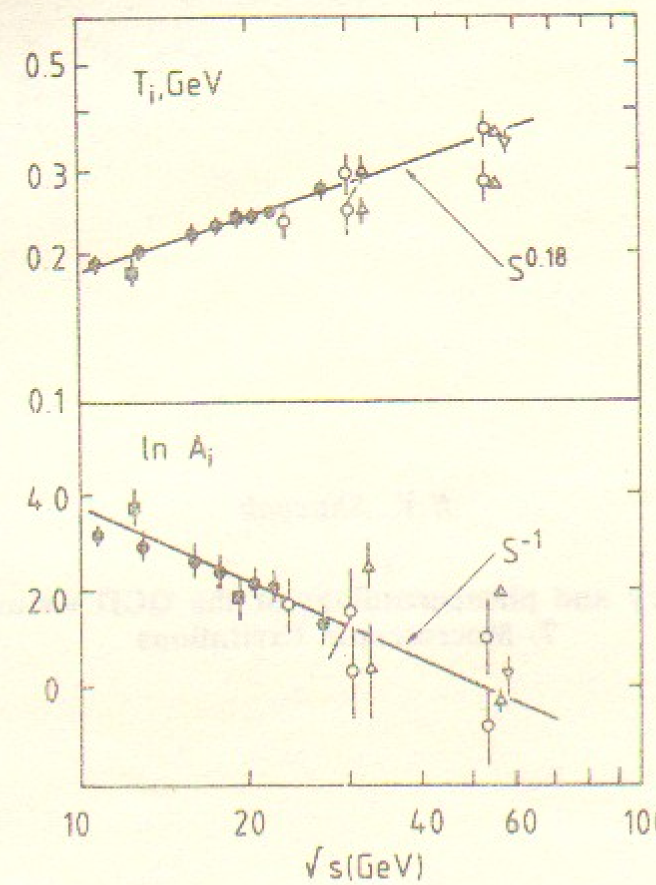

Fig.9. Energy dependence of the «plateau» height T_i and the preexponent factor in the fit A_i. According to the explanation suggested in Ref. [7.59] the former is the intial temperature while the latter is proportional to space-time volume of the intial fierball.

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Theory and phenomenology of the QCD vacuum
7. Macroscopic Excitations

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