УДК 539.121.7

COLD QUARK-GLUON PLASMA
Theoretical and Experimental Perspectives

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The arguments that extremely high-multiplicity hadron interactions at high energies are the source of cold, dense quark-gluon plasma (CQGP) created by the QCD heavy jets are offered. The possibility of calorimetric triggering and measurements of CQGP is considered. The space-time local thermodynamical formalism is adopted for field-theoretical description of such measurements. The valid phenomena in the CQGP are discussed (qualitatively) from theoretical and experimental points of view.

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

Холодная кварк-глюонная плазма. Теоретические и экспериментальные перспективы

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Приводятся аргументы в пользу того, что адронные взаимодействия с экстремально большими множественностями являются источником холодной, плотной кварк-глюонной плазмы (ХКГП), образованной тяжелыми КХД струями. Рассмотрена экспериментальная возможность выделения и измерения ХКГП. Адаптирован локальный в пространстве-времени теоретико-полевой формализм при конечных температурах для описания таких измерений. Возможные явления в ХКГП обсуждаются (качественно) как с теоретической, так и с экспериментальной точек зрения.

Работа выполнена в Лаборатории теоретической физики им. Н.Н.Боголюбова, ОИЯИ.

1. Introduction

It is well known, e.g. [1], that:
— investigation of the transverse energy E_{tr} distribution,
— strange particle creation,
— creation of low and high mass vector mesons;
— creation of direct lepton pairs,
— creation of direct photons

are used as a signal of the quark-gluon plasma (QGP) creation in the high energy interaction. The aim of the talk is to note that:

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— asymptotically high-multiplicity processes are the source of *dense, practically pure, cold* quark-gluon plasma (CQGP). There is no wide discussion of this possibility in the scientific literature. That is why we will start from the very beginning.

The theoretical limit $n \to \infty$ (say $n \geq \bar{n}^2(s)$) will be assumed sometimes since in this limit, as will be seen, the picture of particle creation gets simplified.

This problem hides few questions. First of them:

— *Why the process with $n >> \bar{n}(s)$, where the mean multiplicity $\bar{n}(s)$ introduces the natural scale for n, is the source of CQGP.*

Discussing the *plasma* we assume that it is (i) at least the locally-equilibrium state of (ii) unbounded charges. The first one means that the collective parameters (temperature, density, etc.) are enough to describe the state and the second one is the ordinary condition on the searched state.

This conclusion is not evident since there is also nonperturbative (soft) channel of hadron creation [2], dominated at $n \sim \bar{n}(s)$. Last one describes creation of hadron constituents from vacuum: the kinetic motion of partons leads to increasing, because of confinement phenomena, polarization of the vacuum and to its instability concerning real quark creation [2].

We would like to argue that at the very high multiplicities and energies this effect is negligible, i.e., the polarization of vacuum is small, and cannot shadow both the «direct» particle-creation processes and the collective phenomena in the final-state QGP*. It must be underlined that this is the dynamical, nonkinematical, effect: there is at high energies a wide range of multiplicities, $n_s(s) < n < n_{\max} \sim \sqrt{s}/m$, where $n_s(s) >> \bar{n}(s)$ is the threshold multiplicity, where the vacuum polarization effects are negligible.

We would show the arguments in favour of the perturbative QCD jet mechanism of hadron creation at high multiplicities. It is remarkable also that the number of jets should be minimal (with exponential accuracy at $n \to \infty$) [3]. For example, we valid that practically *total* incident energy is used for «direct» particle production if $n >> \bar{n}(s)$. So, the high-multiplicity processes will allow one to investigate the structure of fundamental Lagrangian since we want to realize the decay of very heavy at high energies initial coloured particle in the «inflational» regime when the vacuum polarization effects are «frozen».

The condition (i) is not so evident. But, considering particles creation as the incident kinetic energies dissipation process, if the multiplicity is the measure of final states entropy, one can assume that the final state is equilibrium at high-multiplicity processes. We would show the quantitative conditions when this is hold.

We hope that:

— *Observation of such (rear) processes will allow to investigate:*

(i) creation of heavy particles (quarks, Higgs bosons, superpartners, etc., as the sources of heavy jets),

(ii) the heavy gluon jets with high E_{tr} creation,

(iii) kinetics of the CQGP formation (noting that the final, preconfinement, state should be near the equilibrium if $n >> \bar{n}(s)$),

*The system of (colour) charges in the polarized vacuum, i.e., when the long-range forces are essential, cannot be considered as the plasma by its ordinary definition.

(iv) collective phenomena (in the dense *cold* pure QGP near the equilibrium).

The second question:

— *How this processes can be triggered.*

We must note that this question is not simple. At $n >> \bar{n}(s)$ the cross sections $\sigma_n(s)$ fall down rapidly and are too small ($< nb$). There is also a problem to adjust the trigger on particle multiplicity.

The high multiplicity experiments imply that at LHC energies there is a possibility to have ten thousand particles in the final state. So, it is natural to think that the strict value of multiplicity n is not important if n is the asymptotically large number. To do first step toward CQGP it is enough to be sure that the transition of «hot» initial state into «cold» final one is examined. For this purpose the ordinary calorimeters can be used for triggering [4]. It is the main physical idea considered below.

We hope that the preparation of such experiment is not a hopeless task and it may be sufficiently informative. This formulation of experiment we will put in the basis of the theory.

Third question:

— *How these processes can be described theoretically.*

In the considered processes we examine (practically total) dissipation of initial-state kinetic energy into particles masses. The theory of dissipative processes has general significance and we concentrated our attention on this important problem.

It must be noted also that the experimental investigation of high-multiplicity processes in a deep asymptotics seems unreal. But considering moderate $n > \bar{n}$ we cannot be sure that the final-state QGP is equilibrium. This leads to necessity to have the theory of dissipation processes with nonequilibrium final state.

It will be offered the economic local-temperatures field-theoretical description based on Wigner function formalism as the mostly close to available experimental layout approach (see also [5]). Our approach can be considered as the natural generalization of ordinary inclusive description since the number of particles, their energies and momentum would not be fixed explicitly (restricting ourselves by calorimetric measurements).

We would discuss particle-creation processes at high energies considering multiplicity n as the characteristics of final state entropy. Under the condition $n >> \bar{n}$ the initial state is very far from equilibrium and we would show that the theory of such processes is simple enough (stationary Markovian) to give definite theoretical predictions. But this leads to the last question:

— *How to define a range of multiplicities, i.e., the value of n_s, where the nonperturbative effects do not play crucial role in the process of particle creation.*

We are not ready to give the answer to this important for experimentators question now, but *we are sure that at sufficiently high energies and sufficiently high multiplicities the hadron interaction processes become hard.*

The discussed problem is located at the cross of a number of today unclear problems. They are the highly nonequilibrium (quantum) thermodynamics, from one hand, and the quantization with conservation laws constraints, from another one. The offered in Conclusion qualitative predictions reflect our today understanding only and should be considered as the illustration to the offered problem.

2. Phenomenology

To build the phenomenology let us introduce the classification of asymptotics over n. It is useful to consider the «big partition function»

$$T(z, s) = \sum_n z^n \sigma_n(s), \quad T(1, s) = \sigma_{tot}(s). \tag{2.1}$$

If we know $T(z, s)$, then $\sigma_n(s)$ is defined by inverse Mellin transformation. This gives (usual in thermodynamics) equation of state:

$$n = z \frac{\partial}{\partial z} \ln T(z, s). \tag{2.2}$$

Solving this equation we can estimate the asymptotics of σ_n:

$$\sigma_n(s) \sim e^{-n \ln \bar{z}(n, s)}, \tag{2.3}$$

where $1 < \bar{z}(n, s) << z_{max}$ is the smallest solution of Eq.(2.2) if the asymptotics over n is considered.

It should be underlined that $\bar{z}(n, s)$ must be an increasing function of n since $\sigma_n(s)$ are positive. Then, it follows from (2.3) that at $n \to \infty$ the solution of (2.2) must tend to singularity located at $z = z_s$ of $T(z, s)$, where $\partial T(z, s)/\partial z$ becomes infinite and the *character of singularity is not important*. So, we must consider three possibilities:

a) $z_s = z_a = 1$, b) $z_s = z_b = \infty$, c) $z_s = z_c$, $1 < z_c < \infty$.

Following Lee and Yang [6] there is no singularities at $0 < z < 1$. Let us consider now the physical content of this classification.

a) $z_s = 1$. It is evident, see (2.1), that in this case at $n \to \infty$

$$0(1) > \sigma_n > 0(e^{-n}), \tag{2.4}$$

i.e., it should decrease but slower than any power of e^{-n}. It is known that the singularity $z_s = 1$ reflects the first-order phase transition [6].

The described mechanism of particle creation assumes that we had prepared a system in the unstable phase and going to another state the system creates particles (this reminds «fate of false vacuum» described in [7]). In hadron physics the initial state may be the QGP and final state should be the hadrons system.

But to have this transition the system should exist. Therefore,

$$\sigma_n \leq 0(e^{-n}) \tag{2.5}$$

since the created hadrons number should be proportional to the number of coloured particles*.

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*The heat capacity of QCD and of hadron system may be different. This can be used for identification of the first order phase transition.

b) $z_s = \infty$. For this case we can put

$$\ln T(z, s) = n_0(s) + \bar{n}(s)(z-1) + O((z-1)^2) \tag{2.6}$$

at $|z-1| << 1$. By definition $n_0(s) = \ln \sigma_{tot}(s)$. The experimental distribution of $\ln T(z, s)$ for various energies shows that the contributions of $O((z-1)^2)$ terms increase with energy [8]. The hadron standard model (SM) assumes that

$$\ln t(z, s) = n_0(s) + \bar{n}_0(s)(z-1)$$

is the Born term in the perturbation series (2.6). There are various interpretations of this series, e.g., the multiperipheral model, the Regge pole model, the heavy color strings model, the QCD multiperipheral models, etc. In all these models $n_0 = a_1 + a_2 \ln s$, $0 \le a_2 << 1$ and $\bar{n}_0(s) = b_1 + b_2 \ln s$, $b_2 > 0$. The second ingredient of hadron SM is the assumption that mean value of created particles transfers momentum $k_{tr} = \text{const}$, i.e., is the energy- and multiplicity independent. It can be shown [8] that under this assumptions $T(z, s)$ is *regular* at finite values of z [8].

We find for this case that $\bar{z}(n, s)$ is the unboundedly increasing function of n. Therefore,

$$\sigma_n < O(e^{-n}) \tag{2.7}$$

for this case.

The SM has a finite range of validity: it is impossible to conserve the condition $k_{tr} = \text{const.}$ if the multiplicities beyond $n \sim \bar{n}^2$ [9]. One can assume that

$$n_s \sim \bar{n}^2. \tag{2.8}$$

c) $1 < z_s < \infty$. Let us assume now that

$$T(z, s) \sim \left(1 - \frac{z-1}{z_s - 1}\right)^{-\gamma}, \quad \gamma > 0. \tag{2.9}$$

Then, using normalization condition, $\partial \ln T(z, s)/\partial z|_{z=1} = \bar{n}_j(s)$ we can find that $z_c(s) = 1 + \gamma/\bar{n}_j(s)$. The singular structure (2.9) is impossible in SM because of condition $k_{tr} = \text{const}$. But if $|z-1| << 1$, we have estimation (2.6). The difference between SM prediction for $T(z, s)$ and (2.9) is seen only at $1 - (z-1)/(z_c - 1) << 1$, i.e., or in asymptotics over n, or in asymptotics over energy.

In considered case $\bar{z} = z_c + 0(\bar{n}_j/n)$ and at high energies $(\bar{n}_j(s) << 1)$

$$\sigma_n \sim e^{-\gamma n/\bar{n}_j} = O(e^{-n}). \tag{2.10}$$

The singular structure (2.9) is familiar for multiplicity distributions in jets. In our terms, if one-jet partition function has the singularity at $z_c^{(1)}(s) = 1 + \gamma \;\; \bar{n}_j(s)$, then the two heavy jets partition function must be singular at

$$z_c^{(2)}(s) = 1 + \frac{\gamma}{\bar{n}_j(s/4)} > z_c^{(1)}(s),$$

because of energy conservation law, and so on. Therefore, at high energies and $n > \bar{n}_j(s)$ the jets number must be minimal with exponential accuracy.

It is evident that the incident energy should be high enough to see this phenomena since the estimation (2.10) predicts the KNO structure.

Comparing (2.7) and (2.10) we can conclude that:

— *At sufficiently high energies the jet mechanism must dominate in asymptotics over n and the number of jets should be minimal; the particle-creation process has a tendency to be stationary Markovian.*

It is the general, practically model-independent, prediction.

So, at high energies there is a wide range of multiplicities where the SM mechanism of hadron creation is negligible. At transition region between «soft» and «hard» channels of hadron creation one can expect the «semihard» processes of mini-jets creation*.

3. Introduction into Formalism

Let us assume that the energies of created particles $\varepsilon_i \leq \varepsilon_0$, where ε_0 is fixed by experiment. For this purpose one can use the calorimeter. Then, using energy conservation law at given ε_0 the number of created particles is bounded from below: $n > \sqrt{s}/\varepsilon_0 \equiv n_{\min}$. With this constraint the integral cross section

$$\sigma_{\varepsilon_0}(s) \sum_{n=n_{\min}} \sigma_n(s)$$

is measured. Choosing $n_{\min} >> \bar{n}$, i.e., $\varepsilon_0 << \sqrt{s}/\bar{n}(s)$, we get into high multiplicity region. There is also a possibility to restore the «differential» cross section $\sim \sigma_n$ calculating the difference $\sigma_{\varepsilon_0}(s) - \sigma_{\varepsilon_0 + \delta\varepsilon_0}(s)$ [4]. It is evident that this method defines n with some accuracy $\pm \Delta n$.

It is not necessary to measure energy of each particle to have $n_{\min} >> \bar{n}$. Indeed, let $\tilde{\varepsilon}_i$ is the energy of i-th group of particles, $\tilde{\varepsilon}_1 + \tilde{\varepsilon}_2 + \ldots + \tilde{\varepsilon}_k = \sqrt{s}$ and let $\tilde{n}_i$ is the number of particles in the group, $\tilde{n}_1 + \tilde{n}_2 + \ldots + \tilde{n}_k = n$. Then, if $\tilde{\varepsilon}_i < \varepsilon_0$, $i = 1,2,\ldots, k$, we have inequality: $k \geq n_{\min}$. Therefore, we get into high multiplicities domain since $n \geq k$ if $\varepsilon_0 << \sqrt{s}/\bar{n}(s)$.

Let us describe now this organization of the experiment. For this purpose we would use the Wigner function formalism in the Carrusers–Zachariasen formulation [10]. For the sake of generality the m into n particles transition probabilities will be considered. This will allow one to include into consideration the heavy ions collisions. If

$$a_{mn}(k; q) = \langle k; m \mid q; n \rangle$$

*The parton lifetime with virtuality $|q|$ is $\sim 1/|q|$ and the time needed for hadron of mass m formation is $\sim 1/m$. Therefore the parton has a time to decay before hadron formation if $|q| >> m$. But this situation is rare because of its kinetic motion.

is the corresponding amplitude, then the m particles interaction cross section with total 4-momentum

$$P = \sum_{i=1}^{m} k_i = \sum_{i=1}^{n} q_i \tag{3.1}$$

is

$$\sigma_{mn}(P) = \int_P \langle k; m | q; n \rangle \langle q; n | k; m \rangle. \tag{3.2}$$

The integration over all k_i and q_i is performed with constraint (3.1).

Considering $d\Omega_m(k)|a_{mn}|$ as the density of states in the Lorentz-invariant phase-space element $d\Omega_m(k)$ of initial state and $d\Omega_n(q)|a_{mn}|$ of final one, the quantity $\sigma_{mn}(P)$ is the density matrix in the energy-momentum representation. We can introduce also the temperature representation considering the temperature as Lagrange multiplier. This is the well-known microcanonical approach in thermodynamics. In the particle physics this idea was explored widely also, see, e.g., [11]. In our case we would introduce two temperatures, for initial state $1/\beta_i$ and for final state $1/\beta_f$ separately since the dissipation processes (transition of kinetic energy into particles masses) are described, $\beta_f >> \beta_i$.

The density matrix in «temperature» representation has the form [12]:

$$\rho(\beta, z) = e^{-\hat{N}(\beta, z; \phi)} R_0(\phi), \tag{3.3}$$

where the operator $\hat{N} = \hat{N}_i + \hat{N}_j^*$ and

$$\hat{N}_{i(f)}(\beta, z; \phi) = \int dr \frac{d^3k}{2\pi^3 2\varepsilon(k)} e^{-\beta_{i(f)}\varepsilon(k)} z_{i(f)}(k, r) \tag{3.4}$$

$$\int dx e^{ikx} \hat{\phi}_{i(f)}(r + z/2) \hat{\phi}_{f(i)}(r - x/2)$$

with $\hat{\phi} \equiv \delta/\delta\phi$ and $\varepsilon(k) = \sqrt{k^2 + m^2}$. Calculating (3.3), the local activities $z_{i(f)}(k, r)$, in analogy with activity z of Sec.2, were introduced. Choosing $z_f(k, r) \sim \Theta(\varepsilon_0 - \varepsilon(k))$, we introduce $n_{\min}$ into formalism.

It is not hard to see that variation of $\rho(\beta, z)$ over $z_{i(f)}$ defines the various cross sections [10]. They define the particles distribution in the phase space (k, r) (we distinguish the initial and final states distributions). The operators $N_{i(f)}(\beta, z; \phi)$ act on the generating functional

$$R_0(\phi) = \int D\Phi_i D\Phi_f e^{iS_0(\Phi_i) - iS_0(\Phi_f)} e^{-iV(\Phi_i + \phi_i) + iV(\Phi_f - \phi_f)}. \tag{3.5}$$

If we put $z_{i(f)} = 1$, $\beta_{i(f)} = \beta$ and calculate $R_0(\phi)$ perturbatively expanding it over V then such defined $\rho(\beta, z = 1)$ coincides [13] with generating functional of Schwinger–Keldysh's real-time finite-temperature field theory [14]. Therefore, the condition $\beta_{i(f)} = \beta$ establishes the isomorphism between our «*S*-matrix» approach and imaginary-time Matsubara theory (of the equilibrium media) [15]. This isomorphism was used in the previous Section.

Let us return now to Eq.(3.3). To find the physical meaning of $\beta_{i(f)}$ we must show the way as they can be measured. If there is nonequilibrium flow, it is hard to invent a thermometer which measures the temperatures of this dissipative processes, i.e., the local in space-time ones. But there is another way — to define the temperatures through equations of state, i.e., as the mean energy of particles. This way is possible in the accelerator experiments where the total energy E is fixed. So, we will define $\beta_{i(f)}$ through equations:

$$E = \frac{\partial}{\partial \beta_{i(f)}} \ln \rho(\beta, z). \tag{3.6}$$

But one cannot find $T(E, z)$ correctly if the assumption that $\beta_{i(f)}$ are «good» quantities is not added, i.e., that the fluctuations near mean values of energy are small (Gaussian).

This assumption is the main problem toward nonequilibrium thermodynamics. The problem in our terms is following: the expansion near $\beta_{i(f)}(E)$ gives asymptotic series with coefficients proportional to the familiar in particle physics inclusive energy spectra:

$$\int \langle\, \varepsilon(k_1), \varepsilon(k_2), \ldots \rangle,$$

where $\langle\,\rangle$ means averaging over interacting fields and integration over k_i is assumed. The careful analysis shows that the fluctuations are small if [15]:

$$\int \langle\, \varepsilon(k_1), \varepsilon(k_2), \ldots \rangle - \prod_i \int \langle\, \varepsilon(k_1)\,\rangle \simeq 0.$$

This factorization property is the well-known Bogoliubov's condition for nonequilibrium thermodynamical systems «shortened» description (in terms of one-particle inclusive cross sections ($\sim \langle\, \varepsilon(k_1)\,\rangle$) only). If $\beta_{i(f)}(E)$ is not the «good« parameter, *all* correlations between created particles are sufficient (one can find discussion of this question in [6]).

It must be noted that this is unique solution of the problem since the considered expansion near $\beta_{i(f)}(E)$ *unavoidably* leads to asymptotic series with zero radii of convergence.

It is hard to expect that $\beta_{i(f)}(E)$ is the good parameter considering the nonequilibrium problem. Nevertheless, there is a possibility to have the factorization averaging over registered ranges of phase space, i.e., in the so-called kinetic phase of the process when the «fast» fluctuations have disappeared and we can consider the long-range fluctuations only.

Then in this domain, with coordinates r and size L, $\beta_{i(f)}(E, r)$ are the good parameters. This is the well-known in nonequilibrium thermodynamics «local equilibrium» hypothesis. We should underline that in our consideration r is the coordinate of *measurement*, i.e., the 4-coordinate where the external particle is measured, and we do not need to divide the interaction region on domains (cells), i.e., will consider r as the calorimeter cells coordinates.

This means that L must be smaller than the typical range of fluctuations. But, on the other hand, L cannot be arbitrary small (since this leads to assumption of *local* factorization property of correlators and, hence, to absence of interactions). This is the natural in quantum theories restriction.

The needed generalization of Wigner function formalism was given in [12]. In this case we must change in (3.4) $\beta \to \beta(r)$ assuming that $\beta_{i(f)}(r)$ and $z_{i(f)}(r)$ are constants on interval L. This prescription adopts Wigner function formalism for the case of high multiplicities. This formalism describes the fluctuations larger than L and averages the fluctuations smaller than L assuming absence, in average, of «non-gaussian» fluctuations. It is the typically «calorimetric» measurement. We will assume that the dimension of calorimeter cells $L < L_{cr}$, where L_{cr} is the dimension of characteristic fluctuations at given n.

In deep asymptotic over n we must have $L_{cr} \to \infty$. The value of particle energies in a cell r is $1/\beta(E, r)$ with exponential accuracy. Last one shows that the offered above experiment with calorimeter as the measuring device for particle energies is sufficiently informative in the very high multiplicities domain.

4. Theoretical Perspectives

Now we would consider the theoretical problem of path integrals (3.5) calculation. Using the equality (unitarity condition):

$$\sum_{m,n} \int \langle k; m \mid q; n \rangle \langle q; n \mid k; m \rangle = \mathrm{Im} \langle \mathrm{vac} \mid \mathrm{vac} \rangle, \tag{4.1}$$

where $\langle \mathrm{vac} \mid \mathrm{vac} \rangle$ is the vacuum-into-vacuum transition amplitude, one can find [17]:

$$\rho(\beta, z) = e^{-i\hat{K}(j,\, e)} \int DM(\Phi) e^{-U(\Phi,\, e)} e^{-N(\beta,\, z;\, \Phi)}, \tag{4.2}$$

where, using Eq.(4.8), $N = N_i + N_j^*$ and

$$N_{i(f)}(\beta, z; \Phi) = \int dr \frac{d^3k}{2\pi^3 2\varepsilon(k)} e^{-\beta_{i(f)}(r)\varepsilon(k)} z_{i(f)}(k, r) \mid \Gamma(k, \Phi) \mid^2 \tag{4.3}$$

with r considered as the *index* of calorimeters cell.

Deriving $N_{i(f)}(\beta, z; \Phi)$, there was used the condition that r is the coordinate of size L cell. With this condition

$$\Gamma(k, \Phi) = \int dx e^{ikx} (\partial_\mu^2 + m^2) \Phi. \tag{4.4}$$

It measures the external fields strength on the remote hypersurface.

The generating functional $R(\beta, z)$ is trivial if $\langle \mid \Gamma(k, \Phi) \mid^2 \rangle = 0$. In this case there is no creation of particles, i.e., there is no measurable asymptotic states (fields). The quantitative condition for this «S-matrix» interpretation of confinement is evident: the *interacting* field $\Phi(x)$ should be equal to zero on the remote hypersurface. In this condition integrating (4.4) by parts we find $\langle \mid \Gamma(k, \Phi) \mid^2 \rangle = 0$.

The expansion over operator

$$\hat{K}(j, e) = 2Re \int dx \frac{\delta}{\delta j(x)} \frac{\delta}{\delta e(x)} \tag{4.5}$$

generates perturbation series and

$$U(\Phi, e) = V(\Phi + e) - V(\Phi - e) - 2Re \int dx e v'(\Phi) \tag{4.6}$$

weights quantum fluctuations. The most important term in (4.2) is the measure

$$DM(\Phi) - \prod_x d\Phi(x)\delta(\partial_\mu^2 \Phi + m^2\Phi + v'(\Phi) - j), \tag{4.7}$$

where $v'(\Phi) \equiv \delta V(\Phi) / \delta\Phi(x)$. So, the equation

$$\partial_\mu^2 \Phi + m^2\Phi + v'(\Phi) = j \tag{4.8}$$

defines the *complete* set of contributions [18].

But Eq.(4.8) gives much more possibilities. Note that l.h.s. of this equation is the sum of classically known forces and the r.h.s. is the quantum force j. Eq.(4.8) establishes the local equality between these forces. This solves the old-standing problem of quantization with constraints.

5. Conclusion

At the end a few words about future steps toward CQGP problem may be useful.

— ***Experimental efforts***

On the today level of understanding we can note that:

A) *At high energies and multiplicities the particle-creation processes would have a tendency to be hard.*

By this reason we can predict that:

a) The mean transverse momentum k_{tr} *should rise with growing multiplicities;*

b) The strange particles mean multiplicity $\bar{n}_{str}$ should rise with total multiplicity n;

c) Must be seen the jet structures;

d) The multiplicity of direct lepton pairs and photons should rise with n.

B) *At high energies and multiplicities the number of jets would have a tendency to be minimal.*

This means that the incident energy would be used effectively for direct particle creation. By this reason:

a) The threshold structures, connected with high-mass fundamental particles, *can* be seen in the (integral) topological cross sections (σ_{ε_0}) σ_n and (or) in the k_{tr} distributions.

C) *At high multiplicities the CQGP becomes equilibrium.*

This can be used for investigation of the thermalization dynamics:

a) The energy spectrum of created particles would have a tendency to be gaussian;

b) The averaged over events energy spectrum and the event-by-event energy distribution in the calorimeter cells would have a tendency to become equal (the ergodic theorem). The divergency from this law may be considered as an indication that not all degrees of freedom were excited in the interaction;

c) The fractal dimensions would change its value in the transition region $n \sim n_s$;

d) The fractal dimensions would be zero for events with $n \gg n_s$.

To realize this programme in the calorimetric experiment the following problems are seen:

(i) The «dead time» of calorimeter should be sufficiently small to use it for «fast» triggering of high-multiplicity events;

(ii) The soft-particle energy measurement accuracy must be sufficiently high;

(iii) The energy loss of small-angle particles should be small (to fix the conservation of total energy).

— Theoretical efforts

We are planning:

a) to extend the perturbative QCD on the high-multiplicity events;

b) to accomodate the used pure phenomenological computer programs for our events generation;

c) to investigate the effectiveness of calorimetric triggering and measurements of the very high multiplicity events.

We hope that this efforts will allow one to formulate the frame of experimental programme.

Beside this we will continue pure theoretical investigation for construction of microscopical theory of hadron interactions based on the formulae (4.2). The work in this direction is in progress now.

Acknowledgement

We would like to thank V.Kadyshevsky for stimulating interest and H.Torosyan for interesting discussions.

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Received on March 12, 1998.