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VERY HIGH MULTIPLICITY PHYSICS

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The Program of Theoretical and Experimental Investigation of Cold Quark-Gluon Plasma

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Abstract

The arguments that the hadrons very high-multiplicity interactions at high energies presents source of could, dense quark-gluon plasma (QGP) are offered. The argument based on classification of asymptotics over multiplicity and the reason of QCD jets dominance is shown. The calorimetric measurements of high multiplicity processes are considered. The corresponding Wigner function formalism is adopted for field-theoretical description of such measurements. The influence of conservation laws on the nonequilibrium flows is discussed.

1 Introduction

It is well known, e.g. [1], that
— investigation of the transfer energy distribution;
— strange particles creation;
— interference effects among identical particles;
— creation of low and high mass vector mesons;
— creation of direct lepton pares;
— creation of direct photons;
— observation of QCD jets
helps find indication of QGP formation. The aim of my talk is to note that
— asymptotically high-multiplicity processes
are the source of *dense, practically pure, cold* quark-gluon plasma (CQGP) formation. I have found only one indication in the scientific literature [2] concerning this possibility. That is why I will start from the very beginning.

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.*
This conclusion is not evident since there are also nonperturbative channels of hadrons creation [3] dominated at $n \sim \bar{n}(s)$. Last ones describe creation of hadrons 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 quarks creation [3]. At the very high multiplicities this effect is negligible and can not shadow the collective phenomena in the *cold* QGP.

I shall start from phenomenological arguments in favor of jet mechanism of hadrons creation at high multiplicities [4]. For this purpose the model independent classification of topological cross sections asymptotics over n will be introduced and the general physical interpretation of classes will be given. I should underline that this conclusion has dynamical reason and is not simple consequence of kinematics.

So, the high-multiplicity processes allows to investigate the structure of fundamental Lagrangian ('unshadowed' by kinetic motions, i.e. we would realize experimentally decay of very hot (at high energies) initial state in the 'inflational' regime, with 'freezed' kinetic degrees of freedom).

The second question:

— *as this processes can be measured.*

I must to note that this question is not simple. At $n >> \bar{n}(s)$ the cross sections $\sigma_n(s)$ falls down rapidly and are too small ($< nb$). There is also a problem to trigger such rear final state.

The high multiplicity experiments imply that at energies of modern accelerators there is hundreds thousand of particles in a final state. It is a hard problem even to count such big numbers. So, the number of particles n can not be considered as a trigger. Moreover, one can think that it is not important have we hundred thousand of particles or hundred thousand plus one. To do first step toward CQGP it is enough to be sure that on experiment the transition of 'hot' initial state into 'cold' final one is examined. For this purpose the ordinary calorimeters can be used [5]. It is the main idea.

The preparation of such experiment is not hopeless task and it may be sufficiently informative. This formulation of experiment we will put in basis of the theory. Theoretically we should shrink the 4-dimension of calorimeter cells up to zero since we do not know *ad hoc* the cells dimension. Then the index of cells i is transformed into the position of particle r. So we come to contradiction with quantum uncertainty principle. This forces to use the Wigner functions formalism [6] and the first question which must be solved is to find a way as this formalism can be adopted for description of our experiment.

Third question:

—*as this processes can be described theoretically.*

In considered processes we examine (practically total) dissipation of initial-state kinetic energy into particles masses. The theory of dissipative processes have general significance from thermodynamical point of view and I would concentrate the attention on this important problem. I want to note also that the experimental investigation of high-multiplicity processes in deep asymptotic seems unreal. But considering moderate $n > \bar{n}$ we can not be sure that the final-state QGP is equilibrium. This leads to necessity to have the theory of dissipation processes with nonequilibrium final state. Nevertheless we will use the language of thermodynamics as the mostly economic formalism, i.e. the way which uses minimal number of parameters (temperature, chemical potential, etc), for description of the large system.

It will be offered the local temperatures field-theoretical description based on Wigner functions formalism as the mostly close to available experimental layout approach. The aim is to show the interconnection between Wigner functions theory and calorimetric experiment (see also [7]).

There is following important observation in nonequilibrium thermodynamics. At the very beginning of this century couple P. and T.Euhrenfest had offered the model to visualize Boltzmann's interpretation of irreversibility phenomena in statistics. The model is extremely simple and fruitful [8]. It considers two boxes with $2N$ numerated balls. Choosing number $l = 1, 2, ..., 2N$ *randomly* one must take the ball with label l from one box and put it to another one. Starting from the 'nonequilibrium' state with all balls in one box it is seen tendency to equalization of balls number in the boxes. So, there is seen irreversible (with insignificant fluctuations!) flow toward preferable, i.e. equilibrium state. This picture shows practical absence[1] of fluctuations in the (nonequilibrium) flow toward a state with maximal entropy.

One can hope that this 'experimental' result reflects a physical reality of nonequilibrium processes with initial state far from equilibrium. I would like in my talk to discuss particles creation processes at high energies from this point of view considering multiplicity n as the characteristics of final state entropy. Under the condition $n >> \bar{n}$ initial state is very far from equilibrium and basing on above described property of nonequilibrium flows one can hope that the theory of such processes is simple enough to give definite theoretical predictions.

If there is not fluctuations in the nonequilibrium flow to high-multiplicity final state one can think that the process is simple Markovian. It is true if there is not long range correlation. Under this special correlations the conservation laws was implied. They are important in dynamics since each conservation law decrease number of degrees of freedom at least on one unite, i.e. it has nonperturbative effect. Moreover, in so-called integrable systems each independent integral of motion (in involution) reduce number of degrees of freedom on two units. In result there is not stochastization in such systems [9], i.e. the nonequilibrium flow is equal to zero. This leads to the last question:
— *as the constraints can be included into formalism).*

The discussed problem is located at the cross of number of today hard problems. They are the highly nonequilibrium (quantum) thermodynamics, from one hand, and the quantization with conservation laws constraints , from another one. In the talk I would like to show some quantitative ideas for this problems solution.

2 Phenomenology

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

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

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

$$n = z\frac{\partial}{\partial z}\ln T(z,s) \qquad (2.1)$$

[1] 'What never? No never! What never? Well, hardly ever.' M.G.Mayer, J.Mayer. *Statistical Mechanics*

Solving this equation we can estimate the asymptotics of σ_n:

$$\sigma_n(s) \sim e^{-n \ln \bar{z}(n,s)}, \tag{2.2}$$

where $1 < \bar{z}(n,s) << z_{max}$ is smallest solution of eq.(2.1).

It follows from (2.2) that at $n \to \infty$ the solution of (2.1) must tend to singularity z_s of $T(z,s)$ and the character of singularity is not important. So, we must consider three possibility:

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

Following to Lee and Yang [10] there is not singularities at $0 < z < 1$. Let us consider now the physical content of this classification.

a). $z_s = 1$. It is evident that

$$\sigma_n \sim e^{-an^{2/3}} > 0(e^{-n}), \quad a > 0, \tag{2.3}$$

i.e. decrease slower then e^{-n}. It is known that the singularity $z_s = 1$ reflects the first order phase transition [10].

To find σ_n for this case we can adopt Langer's analyses [11]. Introducing the temperature $1/\beta$ instead of total energy $\sqrt{s}$ (see Sec.3) we would use the isomorphism with Ising model. For this purpose we divide the volume on cells and if there is particle in the cell we will write (-1). In opposite case (+1). It is the model of lattice gas well described by Ising model. We can regulate the number of down-looking spins, i.e. number of created particles, by the external magnetic field $\mathbf{H}$. Therefore, $z = \exp\{-\beta \mathbf{H}\}$ and $\mathbf{H}$ is the chemical potential.

The described mechanism of particles 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 by Coleman [12]). In hadron physics the initial state may be the QGP and final state may be the hadrons system. Therefore, we must describe the way as the quark-gluon system was prepared.

Following to Lee-Yang's picture of first order phase transition [10](see also [13]) there is not phase transition in a finite system (the partition function can not be singular for finite n_{max}). This means that the multiplicity (and the energy) must be high enough to see described phenomena.

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

$$\ln T(z,s) = n_0(s) + \bar{n}(s)(z-1) + O((z-1)^2) \tag{2.4}$$

at $|z-1| << 1$. By definition $n_0(s) = \ln \sigma_{tot}$. The experimental distribution of $\ln T(z,s)$ for various energies shows that the contributions of $O((z-1)^2)$ terms increase with energy [14] (see Fig.1). The hadrons standard model (SM) assumes that

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

is the Born term in the perturbation series (2.4). There is 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 this models $n_0 = a_1 + a_2 \ln s$, $0 \leq a_2 << 1$ and $\bar{n}(s) = b_1 + b_2 \ln s$, $b_2 > 0$. The second ingredient of hadrons SM is the assumption that mean value of created particles transfers momentum $< k >= const$, i.e. is the energy (and multiplicity) independent. It can be shown that under this assumptions the hadrons SM:

$$\ln T(z,s) = n_0(s) + \sum_n c_n(s)(z-1)^n, \quad c_1 \equiv \bar{n} \tag{2.5}$$

is regular at finite values of z [14].

Inserting (2.5) into (2.1) we find that $\bar{z}(n,s)$ is the increasing function of n. Therefore,

$$\sigma_n < O(e^{-n}). \tag{2.6}$$

But the SM have a finite range of validity: beyond $n \sim \bar{n}^2$ the model must be changed since it is impossible to conserve $< k >= const.$ at higher multiplicities [15].

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

$$T(z,s) \sim (1 - \frac{z-1}{z_c - 1})^{-\gamma}, \quad \gamma > 0. \tag{2.7}$$

Then, using normalization condition, $(\partial T(z,s)/\partial z)|_{z=1} = \bar{n}_j(s)$ we can find that $z_c(s) = 1+\gamma/\bar{n}_j(s)$. The singular structure (2.7) is impossible in SM because of condition $< k >= const.$ But if $|z-1| << 1$ we have estimation (2.4). The difference between SM and c) is seen only at $1-(z-1)/(z_c-1) << 1$, i.e. or in asymptotics over n or in asymptotics over energy. This explains why the asymptotics over n is equivalent of asymptotics over E.

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

$$\sigma_n \sim e^{-\gamma n/\bar{n}_j} = O(e^{-n}). \tag{2.8}$$

—Comparing (2.6) and (2.8) we can conclude that at sufficiently high energies, i.e. if $\bar{n}_j >> \bar{n}$, where $\bar{n}$ is the SM mean multiplicity, the mechanism c) must dominate in asymptotics over n.

It is the general, practically model independent, prediction. It has important from experimental point of view consequence that at high energies there is wide range of multiplicities where the SM mechanism of hadrons creation is negligible. In other words, the CQGP of high multiplicity processes is the dynamical consequence of jets and SM mechanisms. At transition region between 'soft' of SM and 'hard' of jets one can expect the 'semihard' processes of minijets dominance.

The singular structure (2.7) 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 two-jet partition function must be singular at

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

and so on. Therefore, at high energies and $n > \bar{n}_j(s)$ the jets number must be minimal (with exponential accuracy). This means that at $n \to \infty$ the processes of hadrons creation have a tendency to be Markovian (with sharp increase of $< k >$) and only in the last stage of transition (colored plasma) $\to$ (hadrons) the indication on (first order) phase transition may be seen.

One can say that in asymptotics over n we consider the 'inflational' chanel of thermalization which is so fast[2] *that the usual confinement forces becomes 'freezed' and do not play important role in final colored plasma creation.*

Therefore, in asymptotics over n one can expect the transition from 'soft' mechanism of hadrons creation to 'hard' one. The binding forces between colored partons in hard processes are negligible and we can consider this system as the plasma. But previous qualitative analyses allows only to say that the transition occurs.

3 Wigner functions

Let us assume that the energies of created particles $\varepsilon_i \leq \varepsilon_0$, where ε_0 is fixed by experiment. 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 theoretical possibility to restore the quantity $\sim \sigma_n$ calculating the difference $\sigma_{\varepsilon_0}(s) - \sigma_{\varepsilon_0+\delta\varepsilon_0}(s)$ [5] (see Fig.2).

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, $\tilde{\varepsilon}_1 + \tilde{\varepsilon}_2 + ... + \tilde{\varepsilon}_k = \sqrt{s}$ and let $\tilde{n}_i$ is the number of particles in the group, $\tilde{n}_1 + \tilde{n}_2 + ... + \tilde{n}_k = n$[3]*. Then, if $\tilde{\varepsilon}_i < \varepsilon_0$, $i = 1, 2, ..., k$, we have inequality: $k > n_{min}$. Therefor, we get into high multiplicities domain since $n \geq k$, if $\varepsilon_0 << \sqrt{s}/\bar{n}(s)$. We can use the calorimeter demanding that the educed in each cell energy $\tilde{\varepsilon}_i < \varepsilon_0$.*

We will use the Wigner functions formalism in the Carrusers-Zachariasen formulation [6]. For sake of generality the m into n particles transition will be considered. This will allow to include into consideration the heavy ion-ion collisions. If $a_{mn}(k;q)$ is the corresponding amplitude then the m particles interaction cross section with total 4-momentum P is

$$\sigma_{mn}(P) = \frac{1}{n!m!} \int d\Omega_m(k) d\Omega_n(q) \delta(P - \sum_{i=1}^{m} k_i) \delta(P - \sum_{i=1}^{n} q_i) |a_{mn}|^2, \qquad (3.1)$$

[2]The partons life time with virtuality $|q|$ is $\sim 1/|q|$ and the time needed for hadrons of mass m formation is $\sim 1/m$. Therefore the parton have a time to decay before hadrons formation if $|q| >> m$. But this situation is rear because of thermal motion in the initial stage of process is high.

[3]It is assumed that the number of calorimeter cells $K \geq k$.

with Lorenz-invariant phase space element $d\Omega_m(k) = \prod_{i=1}^{m} d^3k_i/(2\pi)^3 2\sqrt{k_i^2+m^2}$. *The amplitudes* a_{mn} *can be computed through the generating functional* $Z(\phi)$ *(the real scalar fields theory is considered for simplicity):*

$$a_{mn}(k_1, k_2, ..., k_m; q_1, q_2, ..., q_n) = \prod_{i=1}^{m} \hat{\phi}(k_i) \prod_{i=1}^{n} \hat{\phi}^*(q_i) Z(\phi),$$

where the annihilation operator is

$$\hat{\phi}(k) = \int dx e^{-ikx} \hat{\phi}(x) \equiv \int dx e^{-ikx} \frac{\delta}{\delta\phi(x)}$$

and $\hat{\phi}^*(q)$ *is the creation operator. The generating functional*

$$Z(\phi) = \int D\Phi e^{iS_0(\Phi) - iV(\Phi+\phi)},$$

where $V(\Phi)$ *describes interactions.*

Considering $d\Omega_m(k)|a_{mn}|$ *as the density of states in the 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 last one as the Lagrange multiplier. This is the well known microcanonical approach of statistics. In the particles physics this idea was explored widely also, see e.g. [16]. 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,* $\beta_f >> \beta_i$.

Let us consider the Fourier transform of δ-functions in (3.1). This introduces two 4-vector α_i *and* α_f*, both conjugate to one 4-vector* P*. We slightly simplify formalism introducing consideration in the CM frame,* $P = (E = \sqrt{s}, \vec{0})$*, and in this case* $\alpha_{i(f)} = (-i\beta_{i(f)}, \vec{0})$.

The density matrix in 'temperature' representation has the form [17]:

$$R(\beta, z) = e^{-\hat{N}(\beta,z;\phi)} R_0(\phi), \tag{3.2}$$

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

$$\hat{N}_{i(f)}(\beta, z; \phi) = \int dr d\Omega_1(k) e^{-\beta_{i(f)}\varepsilon(k)} z_{i(f)}(k, r) \int dx e^{ikx} \hat{\phi}_{i(f)}(r + x/2) \hat{\phi}_{f(i)}(r - x/2). \tag{3.3}$$

with $\hat{\phi} = \delta/\delta\phi$ *and* $\varepsilon(k) = \sqrt{k^2+m^2}$. *Calculating (3.2) the local activities* $z_{i(f)}(k, r)$*, in analogy with activity z of Sec.2, were introduced. It is not hard to see that variation of* $R(\beta, z)$ *over* $z_{i(f)}$ *defines the Wigner functions in Carrusers-Zachariasen [6] representation. They define the particles distribution in the phase space* (k, r) *(we distinguish the initial and final states distributions). The operators* $\hat{N}_{i(f)}(\beta, z; \phi)$ *act on the generating functional*

$$R_0(\phi) = \int D\Phi_i D\Phi_f e^{iS_0(\Phi_i) - iS_0(\Phi_f)} e^{-iV(\Phi_i+\phi_i)+iV(\Phi_f-\phi_f)} \tag{3.4}$$

If we put $z_{i(f)} = 1$, $\beta_{i(f)} = \beta$ and calculate $R_0(\phi)$ perturbatively expanding it over V then such defined $R(\beta)$ coincides [18] with generating functional of Schwinger-Keldysh's real-time finite-temperature field theory [19]. Therefore, the condition $\beta_{i(f)} = \beta$ establish the isomorphysm between our 'S-matrix' approach and imaginary-time Matsubara theory [20]. It is known that last one is the theory of equilibrium states. This isomorphism was used in case a) of previous section.

Let us return now to eq.(3.2). To find the physical meaning $\beta_{i(f)}$ we must show the way as they can be measured. If there is nonequilibrium flow it is hard invent a thermometer (or thermodynamical calorimeter) 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. This way is possible in the accelerator experiments where the total energy E is fixed. So, we will define $\beta_{i(f)}$ through equations:

$$E = \frac{\partial}{\partial \beta_{i(f)}} \ln R(\beta, z), \tag{3.5}$$

i.e. considering $1/\beta_{i(f)}$ as the mean energy of particles in the initial (final) state. But even knowing solutions of this equations one can not find $R(E, z)$ correctly if the assumption that $\beta_{i(f)}$ are 'good' quantities is not added, i.e. that the fluctuations near solutions of eqs.(3.5) 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 over

$$\int \prod \{d\Omega_1(k_i)dr_i\} < \varepsilon(k_1)\varepsilon(k_2)\cdots > |_{(r_1,r_2,\ldots)},$$

where $<>_{()}$ means averaging over fields drown on fixed points of phase space $(k,r)_i$. In other words, the fluctuations near $\beta_{i(f)}(E)$ are defined by value of inclusive spectra familiar in particles physics. Therefore, $\beta_{i(f)}(E)$ are 'good' quantities if this inclusive spectra are small. But this is too strong assumption. More careful analysis shows that it is enough to have the factorization properties [20]:

$$\int \prod \{d\Omega_1(k_i)dr_i\} < \varepsilon(k_1)\varepsilon(k_2)\cdots > |_{(r_1,r_2,\ldots)} - \prod \int d\Omega_1(k_i)dr_i < \varepsilon(k_i) > |_{(r_i)} \sim 0.$$

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

If $\beta_{i(f)}(E)$ is not the 'good' parameter all correlations between created particles are sufficient. And, at the end, discussed factorization property is the well known Bogolyubov's condition of nonequilibrium thermodynamical systems 'shortened' description.

Considering a problem with nonzero nonequilibrium flaw it is hard to expect that $\beta_{i(f)}(E)$ is a good parameter, i.e. that the factorization conditions are hold. Nevertheless there is possibility that in restricted ranges of phase space the mean values of correlators becomes sufficiently small. It is the so called kinetic phase of the process when the memory of initial state was disappeared and the 'fast' fluctuations was disappeared and we can consider the long-range fluctuations only.

Then in this domains, with coordinates r and size L, $\beta_{i(f)}(E,r)$ are the 'good' parameters. This is well known in nonequilibrium thermodynamics the 'local equilibrium' hypothesis. I 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 of QGP on domains (cells), i.e. will consider r as the calorimeter cells coordinates.

This means that L must be smaller then the typical range of fluctuations. But, on other hand, L can not be arbitrary small since this leads to assumption of local factorization property of correlators, i.e. to absence of correlations and, hence, to absence of interactions. This is the natural in quantum theories restriction.

The needed generalization of Wigner functions formalism was given in [17]. In this case we must change in (3.3) $\beta \to \beta(r)$ assuming that $\beta_{i(f)}(r)$ and $z_{i(f)}(r)$ are constants on interval L. This prescription adopts Wigner functions formalism for the case of high multiplicities. This formalism describes the fluctuations larger then L and averages the fluctuations smaller then 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 particles 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 particles energies is sufficiently informative in the very high multiplicities domain.

4 Wigner functions for essentially nonlinear systems

Now we would consider the theoretical problem of path integrals (3.4) calculation. To define the functional measure the ortho-normalizability (i.e. the unitarity) condition will be used. It leads to following representation [21]:

$$R_0(\phi) = e^{-i\hat{K}(j,e)} \int DM(\Phi) e^{-U(\Phi,e)} e^{\int dx (v'(\Phi)+j)\phi}, \tag{4.1}$$

where the expansion over operator

$$\hat{K}(j,e) = 2Re \int dx \frac{\delta}{\delta j(x)} \frac{\delta}{\delta e(x)} \tag{4.2}$$

generates perturbation series and

$$U(\Phi,e) = V(\Phi+e) - V(\Phi-e) - 2Re \int dx e v'(\Phi) \tag{4.3}$$

weights quantum fluctuations. The most important term in (4.1) is the measure

$$DM(\Phi) = \prod_x d\Phi(x) \delta(\partial_\mu^2 \Phi + m^2 \Phi + v'(\Phi) - j) \tag{4.4}$$

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

$$\partial_\mu^2\Phi + m^2\Phi + v'(\Phi) = j \tag{4.5}$$

we will find all contributions[4].

At the very end of calculations one must put $e = j = 0$. *Therefore, eq.(4.4) can be solved expending it over* j. *This shows that (4.1) restores at* $j = 0$ *the usual stationary phase method. Indeed, it can be verified that (4.1) gives usual perturbation theory [21].*

But the eq.(4.5) gives much more possibilities. Note that l.h.s. of this equation is sum of classically known forces and the r.h.s. is the quantum force j. *Eq.(4.5) establish the local equality between this forces. This solves the old standing problem of quantization with constraints: it can be done by field transformations in path integrals since the eq.(4.5) shows the way as* j *must be transformed when the l.h.s. is transformed. Presence of derivatives in (4.5) shows that the quantum force must be transformed in the tangent space of fields. (This explains why the ordinary transformation of path integral (??) is impossible, gives wrong result.)*

We can say that action of operator $e^{-\hat{N}(\beta,z;\phi)}$ *on* $R_0(\phi)$ *maps interacting fields system on measurable states. Let us consider what this gives. Result of action has form:*

$$R(\beta, z) = e^{-i\hat{K}(j,e)} \int DM(\Phi) e^{-U(\Phi,e)} e^{-N(\beta,z;\Phi)}, \tag{4.6}$$

where, using eq.(4.5), $N = N_i + N_f$ *and*

$$N_{i(f)}(\beta, z; \Phi) = \int dr d\Omega_1(k) e^{-\beta_{i(f)}(r)\varepsilon(k)} z_{i(f)}(k, r) |\Gamma(k, \Phi)|^2. \tag{4.7}$$

where r *is considered here as the index of calorimeter cell*[5].

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

$$\Gamma(k, \Phi) = \int dx e^{ikx} (\partial_\mu^2 + m^2)\Phi \tag{4.8}$$

can be considered as the order parameter. Indeed, $\Gamma(k, \Phi)$ *is the element of actions symmetry group since it is linear over field* Φ *and the generating functional* $R(\beta, z)$ *is trivial if* $< |\Gamma(k, \Phi)|^2 >= 0$. *In this case there is not creation of particles, i.e. there is not measurable asymptotic states (fields).*

[4]This means that the unitarity condition is necessary and sufficient for definition of path integral measure in (3.4) [22]

[5]This formulae needs more careful explanation. Instead of (4.7) we must consider

$$N_{i(f)}(\beta, z; \Phi) = \int d\Omega_1(k) \int dr dq \delta(q, L) e^{-\beta_{i(f)}(r)\varepsilon(k)} z_{i(f)}(k, r) \Gamma(k + q, \Phi) \Gamma^*(k - q, \Phi).$$

where L is the scale where $\beta_{i(f)}(r)$ and $z_{i(f)}(k, r)$ can be considered as the constants. In other words, L is the dimension of calorimeter cell. If $L \to \infty$ then $\delta(q, L)$ can be changed on usual δ-function $\delta(q)$ and, therefore, in this limit we will have (4.7). We had considered this limit seance the measurement can not be in contradiction with quantum uncertainty principle.

Consider now the (1+1)-dimensional sin-Gordon model with Lagrange equation:

$$\partial_\mu^2 \Phi + \frac{m^2}{\lambda} \sin(\frac{\lambda}{m}\Phi) = j.$$

At $j = 0$ *this equation has well known soliton solutions* Φ_s *[23] with boundary condition* $\Phi_s(x)|_{|x|=\infty} = 0(mod\ 2\pi\frac{m}{\lambda})$. *It can bee shown [24] that all quantum corrections to solitons contribution in this model equal to zero*[6]. *Then it is easily seen computing integral in (4.8) by parts that* $\Gamma(k, \Phi_s) = 0$. *This result shows that* $\tilde{sl}(2c)$ *symmetry of sin-Gordon model can not be broken and corresponding (polynomial) integrals of motion are conserved.*

5 Conclusion

At the end I wish to say few words about future steps toward CQGP problem.
— Experimental efforts.
They consist in formulation of pure experimental requirements to calorimeters. Here two questions are important: (i) the 'dead time' of calorimeter and (ii) the soft particles energies measurement accuracy. We hope that existing in the modern experiments calorimeters will satisfy our conditions. — Theoretical efforts.
The δ*-likeness of functional measure allows to map a quantum interacting fields system on the principle bundle, where the symmetry constraints are taken into account naturally and the perturbation theory is extremely simple. There is some technical problems. Theirs solution is in progress now. We hope that this will allow to calculate the threshold values of* n *at given energies where the 'hard' processes should dominate.*

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