THE BOSE—EINSTEIN CORRELATIONS FROM THE VIEWPOINT OF QUANTUM FIELD THEORY

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Using a specific version of thermal Quantum Field Theory (QFT), supplemented by operator-field evolution of the Langevin type, we discuss two issues concerning the Bose—Einstein correlations (BEC): the origin of different possible coherent behaviour of the emitting source and the origin of the observed shape of the BEC function $C_2(Q)$. We demonstrate that the previous conjectures in this matter obtained by other approaches are confirmed and have received a complementary explanation.

Introduction

The BEC are recognized since long time as a very important tool providing information about hadronization processes not available otherwise, especially concerning space-time extensions and a coherent or chaotic character of the hadronizing sources. Because the importance of BEC and their present experimental and theoretical status are widely known and well documented (see, for example, [1] and references therein), we shall not repeat it here. Instead we shall proceed to the main point of our interest here, already mentioned above, namely to the discussion of: (*i*) how the possible coherence of a hadronizing system influences the-two body BEC function $C_2(Q)$ [2,3],

$$C_2(Q) = \frac{N_2(k,k')}{N_1(k)N_1(k')} \tag{1}$$

and (*ii*) what is the true origin of the experimentally observed Q-dependence of the $C_2(Q)$ correlation function in the approach used here (out of which the space time information is being deduced, $Q = |k_\mu - k'_\mu| = \sqrt{(k_\mu - k'_\mu)^2}$ with k_μ and k'_μ being the four-momenta of detected particles; in what follows, we assume for simplicity that all produced particles are bosons). In literature, one finds that, in some approaches using quantum statistical methods [2],

$$C_2(Q) = 1 + 2p(1-p)\sqrt{\Omega(q)} + p^2\Omega(q), \tag{2}$$

whereas in other approaches [3]

$$C_2(Q) = 1 + \lambda\Omega(Qr) \tag{3}$$

(actually, this is the most frequently used form). In both cases, the parameters p and λ are called *coherence* parameters defining the degree of coherence of the hadronizing source (for a purely coherent source, $p = \lambda = 0$ and there is no BEC, for $p = \lambda = 1$, i.e., in the purely chaotic case, both equations coincide). Although in [2, 3] they are operationally expressed in the same way, i.e., $p = \lambda = \langle N_{\rm chaotic}\rangle / \langle N_{\rm total}\rangle$, we shall formally differentiate between them when using the corresponding expressions for $C_2(Q)$ because, as will be shown later, the concepts of coherence they correspond to are different in each case. The choices of $\Omega(Q)$ discussed in literature vary between [4–6] (here $q = rQ$ with $r = |r_\mu| = \sqrt{r_\mu r_\mu}$ being a 4-vector such that $\sqrt{(r_\mu)^2}$ has dimension of length and the product $Qr = Q_\mu r_\mu = q$ is dimensionless):

- Gaussian: $\Omega(q) = \exp\left(-Q^2r^2\right)$;
- exponential: $\Omega(q) = \exp(-Qr)$;
- Lorentzian: $\Omega(q) = 1/\left(1+Qr\right)^2$;
- given by the Bessel function [5]: $\Omega(q) = \left[J_1(Qr)/(Qr)\right]^2$.

The most frequently used forms are the Gaussian and exponential ones [1]. The above questions can be therefore rephrased in the following way: (*i*) why are forms of $C_2(Q)$ in Eqs. (2) and (3) different and are parameters p and λ referring to the same quantity and (*ii*) what stays behind a specific choice of the form of the $\Omega(Q)$ function as listed above.

To make our point more clear, we shall work here directly in phase space (as, for example, in [7]), and no space-time considerations will be used (contrary to the majority of works on BEC [1–3]). As our working tool, we shall choose some specific (thermal) version of Quantum Field Theory (QFT) supplemented by the operator-field evolution of the Langevin type proposed recently [8,9]. We shall demonstrate that: (i) the origin of differences in $C_2(Q)$ in Eqs. (2) and (3) lies in different ways of introducing the concept of coherence in both approaches, i.e., p and λ referred to different concepts of coherence in each case; (ii) in order to obtain a given (experimentally observed) shape of the BEC correlation function $C_2(Q)$ (i.e., $\Omega(Q)$), one has to account somehow for the *finiteness* of the space-time region of the particle production (i.e., of the hadronizing *source*). In QFT approach used here, it is particularly clearly seen and is connected with the necessity of smearing out some generalized functions (delta functions: $\delta(Q_\mu = k_\mu - k'_\mu)$) appearing in the definition of thermal averages of some operators occurring here. The freedom in using different types of smearing functions to perform such a procedure allows us to account for all possible different shapes of hadronizing sources apparently observed by experiment. (Actually, a careful inspection of all previous approaches to BEC using QFT, cf., for example, [11], shows that this was always the procedure used, though never expressed so explicitly as is done here.).

1. Description of a Hadronizing Source

Let us recapitulate now the main points of our approach (for details, see [8]). The collision process produces usually a large number of particles, out of which we select one (we assume for simplicity that we are dealing only with identical bosons) and describe it by the operator $b(\vec{k},t)$ (the notation is the usual one: $b(\vec{k},t)$ is an annihilation operator, $\vec{k}$ is 3-momentum and t is a real time). The rest of particles is then assumed to form a kind of heat bath, which remains in equilibrium characterized by a temperature $T = 1/\beta$ (which will be one of our parameters). All averages $\langle(\dots)\rangle$ are therefore thermal averages of the type:

$$\langle(\dots)\rangle = \mathrm{Tr}\left[(\dots)e^{-\beta H}\right]/\mathrm{Tr}\left(e^{-\beta H}\right). \tag{4}$$

However, we shall also allow for some external (to the above heat bath) influence acting on our system. Therefore we shall represent the operator $b(\vec{k},t)$ as that consisting of a part corresponding to the action of the heat bath, $a(\vec{k},t)$, and also of a part describing the action of these external factors, $R(\vec{k},t)$:

$$b(\vec{k},t) = a(\vec{k},t) + R(\vec{k},t). \tag{5}$$

The time evolution of such a system is then assumed to be given by the Langevin equation [9]

$$i\partial_t b(\vec{k},t) = F(\vec{k},t) - A(\vec{k},t) + P \tag{6}$$

(and a similar conjugate equation for $b^+(\vec{k},t)$). These equations are supposed to model all aspects of the hadronization process (notice their similarity to the equations describing the motion of a Brown particle in some external field [10]). The meaning of different terms appearing here is the following:

(i) The combination $F(\vec{k},t) - A(\vec{k},t)$ represents the so-called *Langevin force* and is therefore responsible for the internal dynamics of hadronization in the following manner:

- A is related to stochastic dissipative forces and is given by [8, 9]

$$A(\vec{k},t) = \int\limits_{-\infty}^{+\infty} d\tau K(\vec{k},t-\tau)b(\vec{k},\tau), \tag{7}$$

  with the operator $K(\vec{k},t)$ being a random evolution field operator describing the random noise and satisfying the usual correlation-fluctuation relation for the Gaussian noise: $\langle K^+(\vec{k},t)K(\vec{k}',t)\rangle = 2\sqrt{\pi\rho}\kappa\delta(\vec{k}-\vec{k}')$ (κ and ρ are parameters describing the effect caused by this noise on the particle evolution in thermal environment [9]).

- The operator $F(\vec{k},t)$ describes the influence of the heat bath,

$$F(\vec{k},t) = \int\limits_{-\infty}^{+\infty} \frac{d\omega}{2\pi}\psi(k_\mu)\hat{c}(k_\mu)e^{-i\omega t}. \tag{8}$$

(ii) Our heat bath is represented by an ensemble of damped oscillators, each is described by the operator $\hat{c}(k_\mu)$ such that $\left[\hat{c}(k_\mu),\hat{c}^+(k'_\mu)\right] = \delta^4(k_\mu - k'_\mu)$, and characterized by some function $\psi(k_\mu)$, which is subjected to a kind of normalization involving also dissipative forces represented by the Fourier transformed operator $\tilde{K}(k_\mu)$, namely:

$$\int\limits_{-\infty}^{+\infty} \frac{d\omega}{2\pi}\left[\frac{\psi(k_\mu)}{\tilde{K}(k_\mu)-\omega}\right]^2 = 1. \tag{9}$$

(iii) Finally, the constant term P (representing *the external source* term in the Langevin equation) denotes the possible influence of some external force (assumed here to be constant in time). This force would result, for

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example, in a strong ordering of phases, leading therefore to the coherence effect in the sense discussed in [2].

Out of many details (for which we refer to [8]), what is important in our case is the fact that the 2-particle correlation function for like-charge particles, as defined in (1), is given in such a form ($(k_\mu = (\omega = k^0, k_j)$):

$$C_2(Q) = \xi(N) \frac{\tilde{f}(k_\mu, k'_\mu)}{\tilde{f}(k_\mu)\tilde{f}(k'_\mu)} = \\ = \xi(N)\left[1 + D(k_\mu, k'_\mu)\right], \quad (10)$$

where

$$\tilde{f}(k_\mu, k'_\mu) = \langle \tilde{b}^+(k_\mu)\tilde{b}^+(k'_\mu)\tilde{b}(k_\mu)\tilde{b}(k'_\mu)\rangle, \\ \tilde{f}(k_\mu) = \langle \tilde{b}^+(k_\mu)\tilde{b}(k_\mu)\rangle \quad (11)$$

are the corresponding thermal statistical averages (in which temperature T enters as a parameter) with $\tilde{b}(k_\mu) = \tilde{a}(k_\mu) + \tilde{R}(k_\mu)$ being the corresponding Fourier transformed stationary solution of Eq. (6). As shown in [8] (notice that the operators $\tilde{R}(k_\mu)$ by definition commute with themselves and with any other operator considered here):

$$\tilde{f}(k_\mu, k'_\mu) = \tilde{f}(k_\mu)\tilde{f}(k'_\mu) + \\ +\langle \tilde{a}^+(k_\mu)\tilde{a}(k'_\mu)\rangle\langle \tilde{a}^+(k'_\mu)\tilde{a}(k_\mu)\rangle + \\ +\langle \tilde{a}^+(k_\mu)\tilde{a}(k'_\mu)\rangle\tilde{R}^+(k'_\mu)\tilde{R}(k_\mu) + \\ +\langle \tilde{a}^+(k'_\mu)\tilde{a}(k_\mu)\rangle\tilde{R}^+(k_\mu)\tilde{R}(k'_\mu), \quad (12)$$

$$\tilde{f}(k_\mu) = \langle \tilde{a}^+(k_\mu)\tilde{a}(k_\mu)\rangle + |\tilde{R}(k_\mu)|^2. \quad (13)$$

This defines $D(k_\mu, k'_\mu)$ in (10) in terms of the operators $\tilde{a}(k_\mu)$ and $\tilde{R}(k_\mu)$, which are equal in our case to:

$$\tilde{a}(k_\mu) = \frac{\tilde{F}(k_\mu)}{\tilde{K}(k_\mu) - \omega} \quad \text{and} \quad \tilde{R}(k_\mu) = \frac{P}{\tilde{K}(k_\mu) - \omega}. \quad (14)$$

The factor ξ depending on the multiplicity N is, in our case, equal to $\xi(N) = \langle N(N-1)\rangle/\langle N\rangle^2$. This means therefore that the correlation function $C_2(Q)$, as defined by Eq. (10), is essentially given in terms of P and the two following thermal averages for the $F(\vec{k}, t)$ operators:

$$\langle F^+(\vec{k}, t)F(\vec{k}', t')\rangle = \delta^3(\vec{k} - \vec{k}') \times \\ \times \int \frac{d\omega}{2\pi}\, |\psi|^2\, n(\omega) e^{+i\omega(t-t')}, \\ \langle F(\vec{k}, t)F^+(\vec{k}', t')\rangle = \delta^3(\vec{k} - \vec{k}') \times \\ \times \int \frac{d\omega}{2\pi}\, |\psi|^2\, [1 + n(\omega)] e^{-i\omega(t-t')} \quad (15)$$

where $n(\omega) = \{\exp[(\omega - \mu)\beta] - 1\}^{-1}$ is the number of (by assumption - only bosonic in our case) damped oscillators of energy ω in our reservoir characterized by parameters μ (chemical potential) and inverse temperature $\beta = 1/T$ (both being free parameters). The origin of these parameters, the temperature $\beta = 1/T$ and chemical potential μ, is the Kubo—Martin—Schwinger condition that

$$\langle a(\vec{k}', t')a^+(\vec{k}, t)\rangle = \langle a^+(\vec{k}, t)a(k', t - i\beta)\rangle \exp(-\beta\mu), \quad (16)$$

(see [8, 9]). This form of averages presented in (15) reflects the corresponding averages for the $\hat{c}(k_\mu)$ operators, namely that

$$\langle \hat{c}^+(k_\mu)\hat{c}(k'_\mu)\rangle = \delta^4(k_\mu - k'_\mu)n(\omega), \\ \langle \hat{c}(k_\mu)\hat{c}^+(k'_\mu)\rangle = \delta^4(k_\mu - k'_\mu)[1 + n(\omega)]. \quad (17)$$

Notice that, with only delta functions present in (15), one would have a situation in which our hadronizing system would be described by some kind of *white noise* only. The integrals multiplying these delta functions and depending on (a) momentum characteristic of our heat bath $\psi(k_\mu)$ (representing in our case, by definition, the hadronizing system) and (b) assumed bosonic statistics of produced secondaries resulting in the factors $n(\omega)$ and $1+n(\omega)$, respectively, bring the description of our system closer to reality.

2. Results

It is straightforward to realize that the existence of BEC, i.e., the fact that $C_2(Q) > 1$, is strictly connected with nonzero values of the thermal averages (15). However, in the form presented there, they differ from zero *only at one point*, namely for $Q = 0$ (i.e., for $k_\mu = k'_\mu$). Actually, this is the price one pays for the QFT assumptions tacitly made here, namely for the *infinite* spatial extension and for the *uniformity* of our reservoir. But we know from the experiment [1] that $C_2(Q)$ reaches its maximum at $Q = 0$ and falls down towards its asymptotic value of $C_2 = 1$ at large Q (actually already at $Q \sim 1$ GeV/c). To reproduce the same behaviour by means of our approach here, one has to replace delta functions in Eq. (18) by functions with supports larger than those limited to a one point only. This means that such functions should not be infinite at $Q_\mu = k_\mu - k'_\mu = 0$ but remain more or less sharply peaked at this point, otherwise remaining finite and falling to zero at small, but finite, values of $|Q_\mu|$ (actually the same as those at which $C_2(Q)$ reaches unity):

$$\delta(k_\mu - k'_\mu) \Longrightarrow \Omega_0\sqrt{\Omega(q = Qr)}. \quad (18)$$

Here, Ω_0 has the same dimension as the δ function (actually, it is nothing else but a 4-dimensional volume restricting the space-time region of particle production) and $\Omega(q)$ is a dimensionless smearing function which contains the q-dependence we shall be interested in here. In this way, we are tacitly introducing a new parameter (mentioned already in Introduction), r_μ, a 4-vector such that $\sqrt{(r_\mu)^2}$ has dimension of length and which makes the product $Qr = Q_\mu r_\mu = q$ dimensionless. This defines the region of *nonvanishing* density of oscillators $\hat{c}$, which we shall *identify* with the space-time extensions of the hadronizing source. Expression (18) has to be understood in a symbolic sense, i.e., that $\Omega(Qr)$ is a function which becomes *strictly* a δ-function in the limit $r \to \infty$. Making such a replacement in Eq. (15), one must also decide how to accordingly adjust $n(\omega)$ occurring there because now, in general, $\omega \neq \omega'$. In what follows, we shall simply replace $n(\omega) \to n(\bar{\omega})$ with $\bar{\omega} = (\omega + \omega')/2$ (which, for classical particles, would mean that $n(\omega) \to \sqrt{n(\omega)n(\omega')}$).

In such a way, r becomes a new (and from the phenomenological point of view also the most important) parameter entering here together with the whole function $\Omega(Qr)$, to be deduced from comparison with experimental data (one should notice that the opposite line of reasoning has been used in [12] where, at first, a kind of our $\Omega(q)$ function was constructed for a finite source function and it was then demonstrated that, in the limit of infinite homogenous source, one ends with a delta function). With such a replacement, one now has

$$D(k_\mu, k'_\mu) = \frac{\sqrt{\tilde{\Omega}(q)}}{(1+\alpha)(1+\alpha')}\left[\sqrt{\tilde{\Omega}(q)} + 2\sqrt{\alpha\alpha'}\right], \quad (19)$$

where

$$\tilde{\Omega}(q) = \gamma\Omega(q), \quad \gamma = \frac{n^2(\bar{\omega})}{n(\omega)n(\omega')}, \quad \alpha \propto \frac{P^2}{|\psi(k_\mu)|^2 n(\omega)}, \quad (20)$$

with $n(\omega)$ to be the same as defined above.

Another very important parameter entering (19) is α, which, first of all, reflects the action of external force P present in the evolution equation (6). This action is combined here (in a multiplicative way) with information on both the momentum dependence of the reservoir (via $|\psi(k_\mu)|^2$) and on the single particle distributions of the produced particles [via $n(\omega = \mu_T \cosh y)$, where μ_T and y are, respectively, the transverse mass and rapidity]. Parameter α summarizes therefore our knowledge of other characteristics than space-time ones of the hadronizing source (given by $\Omega(q)$ introduced above). Notice that $\alpha > 0$, only when $P \neq 0$. Actually, for $\alpha = 0$ one has

$$1 < C_2(Q) < 1 + \gamma\Omega(Qr), \quad (21)$$

i.e., it is contained between the limits corresponding to very large (lower limit) and very small (upper limit) values of P. Because of this, α plays the role of a *coherence* parameter [1, 2]. For $\gamma \simeq 1$, neglecting the possible energy-momentum dependence of α and assuming that $\alpha' = \alpha$, one gets the expression

$$C_2(Q) = 1 + \frac{2\alpha}{(1+\alpha)^2}\sqrt{\Omega(q)} + \frac{1}{(1+\alpha)^2}\Omega(q), \quad (22)$$

which is formally *identical* with Eq. (2) obtained in [2] by means of the QS approach. It has precisely the same form, including two Q-dependent terms containing the information on the shape of the source, one being the square of the other, each multiplied by some combination of the *chaoticity* parameter $p = 1/(1+\alpha)$ (however, in [2], p is defined as the ratio of the mean multiplicity of particles produced by the so-called *chaotic* component of the source to the mean total multiplicity, $p = \langle N_{\rm ch}\rangle/\langle N\rangle$). In fact, because in general $\alpha \neq \alpha'$ (due to the fact that $\omega \neq \omega'$ and therefore the number of states, identified here with the number of particles with given energy, $n(\omega)$, is also different), one should rather use the general form (10) for C_2 with details given by (19) and (20) and with α depending on such characteristics of the production process as temperature T and chemical potential μ occurring in the definition of $n(\omega)$.

Notice that Eq. (22) differs from the usual empirical parametrization of $C_2(Q)$ [1] as given by Eq.(3) in which $0 < \lambda < 1$ is a free parameter adjusting the observed value of $C_2(Q=0)$, which is customary called "incoherence" and with $\Omega(Qr)$ represented usually as Gaussian) should be replaced by Eq. (2) with $\alpha = (1 - \sqrt{\lambda})/\sqrt{\lambda}$. Recently Eq. (3) has found a strong theoretical support expressed in detail in [3] (where λ is given by the same ratio of multiplicities as p above). The natural question arises: which of the two formulas presented here is correct? The answer we propose is: both are right in their own way. This is because each of them is based on different ways of defining the coherence of a source. In [3], one uses the notion of coherently and chaotically produced particles or, in other words, one divides a hadronizing source into coherent and chaotic subsources. In [2], one introduces instead the notion of partially coherent fields representing produced particles, i.e., one has only one source which produces partially coherent fields. Our approach is similar, as we describe our particle by the operator $b(\vec{k}, t)$, which consists of two parts, cf. Eq. (5), one of which depends on the external static force P. The action of this force is to *order the phases* of particles in our source (represented by the heat bath). The strength of this ordering depends on the value of the external force P. In any case, for $P \neq 0$, it demonstrates itself as a *partial coherence.*

Some comments are necessary at this point. Notice that it is of the same type as that considered in [2]. When comparing with [3], one should notice that although our operators $\tilde{a}(k_\mu)$ and $\tilde{R}(k_\mu)$ look similar to operators defined in Eqs. (4) and (5) of [3] they differ in the following. Our $R(k_\mu)$ describes essentially the action of *a constant* force P and as such it commutes with all other operators (including themselves). So it only introduces a partial ordering of the phases of particles decreasing the C_2 correlation function, i.e., acting as a coherent component, albeit we do not have coherent particles as such. It is also seen when realizing that in Eq. (13) the two last terms contain only one pair of operators a. This, in the language of [3], translates to only one Wigner function, $f_{\rm ch}$, to be present here. The operators R cannot form the second Wigner function ($f_{\rm coh}$ in [3]). This is the technical origin of the three terms present in (2) (and in [2]) in comparison to two terms in (3) and obtained in [3].

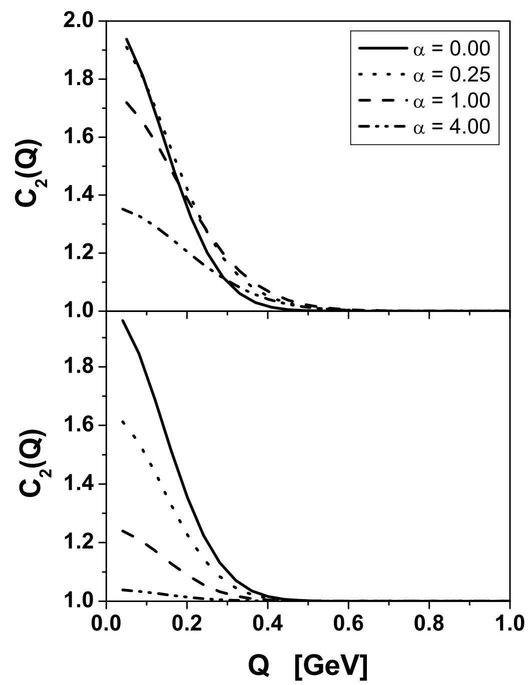
Let us return to the problem of the Q-dependence of BEC. One more remark is here. The problem with the $\delta(k_\mu - k'_\mu)$ function encountered in two-particle distributions does not exist in the single-particle distributions, which are in our case given by Eq. (13) and which can be written as $\tilde{f}(k_\mu) \propto \langle \tilde{a}^+(k_\mu)\tilde{a}(k_\mu)\rangle + |\tilde{R}(k_\mu)|^2 \sim (1+\alpha)\langle \tilde{a}^+(k_\mu)\tilde{a}(k_\mu)\rangle$ (it is normalized to the mean multiplicity: $\int d^4k\, \tilde{f}(k_\mu) = \langle N\rangle$). To be more precise

$$\tilde{f}(k_\mu) = (1+\alpha)\Xi(k_\mu, k_\mu), \tag{23}$$

where $\Xi(k_\mu, k_\mu)$ is the one-particle distribution function for the "free" (undistorted) operator $\tilde{a}(k_\mu)$ equal to

$$\Xi(k_\mu, k_\mu) = \Omega_0 \left| \frac{\psi(k_\mu)}{\tilde{K}(k_\mu) - \omega} \right|^2 n(\omega). \tag{24}$$

Notice that the actual shape of $\tilde{f}(k_\mu)$ is dictated both by $n(\omega) = n(\omega; T, \mu)$ (calculated for fixed temperature T and chemical potential μ at energy $\hat{\omega}$ as given by the Fourier transform of the random field operator $\tilde{K}$ and by shape of the reservoir in the momentum space provided by $\psi(k_\mu)$) and by external force P in parameter α. They are both unknown, but because these details do not enter the BEC function $C_2(Q)$, we shall not pursue this



Shapes of $C_2(Q)$ as given by Eq. (2) — upper panel and for the truncated version of (2) (without the middle term) corresponding to Eq. (3) — lower panel. Gaussian shape of $\Omega(q)$ was used in both cases

problem further. What is important for us at the moment is that both the coherent and the incoherent parts of the source have the same energy-momentum dependence (whereas in other approaches mentioned here they were usually assumed to be different). On the other hand, it is clear from (23) that $\langle N\rangle = \langle N_{\rm ch}\rangle + \langle N_{\rm coh}\rangle$ (where $\langle N_{\rm ch}\rangle$ and $\langle N_{\rm coh}\rangle$ denote the multiplicities of particles produced chaotically and coherently, respectively) therefore justifying the definition of chaoticity p mentioned above.

Figure shows in detail (using Gaussian shape of $\Omega(q)$ function) the dependence of $C_2(Q)$ on different values of $\alpha = 0,\ 0.25,\ 1,\ 4$ (again, used in the same approximate way as before and corresponding to $p = 1,\ 0.8,\ 0.5,\ 0.2$) and compare it to the case where the second term in Eq. (2) is neglected, as is the case in the majority of phenomenological fits to data.

3. Summary and Conclusions

To summarize: using a specific version of QFT supplemented by the Langevin evolution equation (6) to describe the hadronization process [8, 9], we have derived the usual BEC correlation function in the form explicitly showing the origin of both the so-called coherence (and how it influences the structure of BEC) and the Q-dependence of BEC represented by the correlation function $C_2(Q)$. The dynamical source of coherence is identified in our case with the existence of a constant external term P in the Langevin equation. Its influence turns out to be identical with the one obtained before in the QS approach [2] and is described by Eq. (19). Its action is to order the phases of produced secondaries. Therefore for $P \to \infty$, we have all phases aligned in the same way and $C_2(Q) = 1$. This is because both here and in [2] the coherence has already been introduced on the level of a hadronizing source, as a property of fields (in [2]) or operators describing produced particles. Dividing instead the hadronizing source itself into coherent and chaotic subsources leads to results obtained in [3] and given by Eq. (3). The controversy between the results given by [2] and [3] is therefore explained: both approaches are right, one should only remember that they use different descriptions of the notion of coherence. It is therefore up to the experiment to decide which proposition is followed by nature: the simpler formula (3) or rather the more involved (10) together with (19). From Fig. 1, one can see that differences between both forms are clearly visible, especially for larger values of coherence α, i.e., for the lower chaoticity parameter p.

From our presentation, it is also clear that the form of C_2 reflects distributions of the space-time separation between the two observed particles rather than the distribution of their separate production points (cf., for example, [13], where it is advocated that it is in fact a Fourier transform of the two-particle density profile of a hadronizing source, $\rho(r_1, r_2) = \rho(r_1 - r_2)$, without approximating it by the product of single-particle densities, as in [1]).

Finally, we would like to stress that our discussion is so far limited to only a single type of secondaries being produced. It is also aimed at a description of hadronization understood as a kinetic freeze-out in some more detailed approaches. So far we were not interested in the other (highly model dependent) details of the particle production process. This is enough to obtain our general goals, i.e., to explain the possible dynamical origin of coherence in BEC and the origin of the specific shape of the correlation $C_2(Q)$ functions as seen from the QFT perspective. We close with the suggestion that both sources of coherence, that presented here and in [2] and that investigated in [3], should be considered together. The most general situation would be then a hadronizing source composed with a number of subsources, each with different internal (discussed here) degree of coherence. For only one subsource present, we would have situation described here whereas for a number of subsources, each being either totally coherent or totally chaotic, the description offered by [3] would automatically emerge. It is up to experimental data to decide and Fig. 1 tells us that it is not impossible task (at least in principle).

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БОЗЕ-ЕЙНШТЕЙНІВСЬКІ КОРЕЛЯЦІЇ З ТОЧКИ ЗОРУ КВАНТОВОЇ ТЕОРІЇ ПОЛЯ

Г. Козлов, О. Утюж, Г. Вілк

Р е з ю м е

Використовуючи спеціальну версію теплової квантової теорії поля, доповнену операторно-польовою еволюцією ланжевенівського типу, ми обговорюємо два аспекти бозе-ейнштейнівських кореляцій: походження іншої можливої когерентної поведінки джерела випромінювання і походження форми бозе-ейнштейнівської кореляційної функції $C_2(Q)$, що вивчається. Ми показуємо, що попередні припущення в цій області, одержані в рамках інших підходів, підтверджуються і отримують додаткове пояснення.

БОЗЕ-ЭЙНШТЕЙНОВСКИЕ КОРРЕЛЯЦИИ С ТОЧКИ ЗРЕНИЯ КВАНТОВОЙ ТЕОРИИ ПОЛЯ

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Р е з ю м е

Используя специальную версию тепловой квантовой теории поля, дополненную операторно-полевой эволюцией ланжевеновского типа, мы обсуждаем два аспекта бозе-эйнштейновских корреляций: происхождение другого возможного когерентного поведения излучающего источника и происхождение наблюдаемой формы бозе-эйнштейновской корреляционной функции $C_2(Q)$. Мы показываем, что предыдущие предположения в этой области, полученные в рамках других подходов, подтверждаются и получают дополнительное объяснение.