Higher Order Bose-Einstein Correlations for Systems with Large Halo

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Abstract

A generalization of the core/halo model to the Bose-Einstein correlation function of any number of particles is presented. In particular, a simple prediction is obtained for the intercept parameter of the \(n\)-particle Bose-Einstein correlation functions for arbitrary large values of \(n\).

Introduction. Multi-particle correlations are studied extensively in high energy collisions, with the aim of learning more about possible phase transitions or self-similar fluctuations. Especially, the short-range part of the two-particle correlation function, is thought to carry information about the space-time structure based on a quantum-statistical effect discovered in refs. [1, 2]. One of the most important driving force behind correlation studies in high energy physics is the possibility to measure space-time dimensions on the \(10^{-15}\) m and \(10^{-23}\) sec scale. See refs. [3, 4] for recent reviews on high energy heavy ion physics, for summaries on correlation studies in this field see refs. [5, 6, 7, 8].

Multi-particle correlations are becoming experimentally well determinable in present and planned heavy ion reactions at at Brookhaven AGS, at CERN SPS, at the Relativistic Heavy Ion Collider (RHIC) and at CERN LHC both due to the very large number of the produced particles and due to the dramatic increase of the data quality accessible already for multidimensional correlation studies, ref. [9].

There has been a lot of development recently in the theory of multi-particle Bose-Einstein correlations, form partial coherence to multi-particle wave-packet models and event generators with multi-particle symmetrization, refs. [10, 11, 12, 13, 14, 15, 16, 17]. The purpose of the present Letter is to study the structure of multi-particle Bose-Einstein correlations in one special case, when the emission function can be well separated to a

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core and a halo part. Such a scenario is referred to as the core/halo model [18]. The formal study of the core/halo type of models first started with numerical simulations, [19, 20, 21, 22, 23, 24]. The essential ideas necessary to formulate the core/halo model were also published in ref. [25], see refs. [26, 27] for further details.

In refs. [28, 29, 30] an Edgeworth expansion method was proposed to characterize the non-Gaussian features of the Bose-Einstein correlation functions as caused by interference of decay product pions from long-lived resonances. However, the data indicated surprisingly regular, simple Gaussian correlation functions when the analysis was performed in two or three dimensions [31, 32]. This resulted [33] in the re-formulation and simplification of the theoretical aspects of the resonance-decay effects on the Bose-Einstein correlation functions in a particularly transparent form, with an emphasis on the interpretation and possible utilization of the momentum dependent intercept parameter of the two particle Bose-Einstein correlation function in ref. [18]. We present herewith the generalization of this core/halo model to the \( n \)-particle case, utilizing the Wigner-function formalism.

In this formalism, the one-boson emission is characterized by the emission function, \( S(x, p) \). Here \( x = (t, \mathbf{r}) \) and \( p = (E, \mathbf{p}) \) denote four-vectors in space-time and in momentum space, for particles with mass \( m = \sqrt{E^2 - \mathbf{p}^2} \). An auxiliary quantity is

\[
\tilde{S}(\Delta k, K) = \int d^4x \ S(x, K) \exp(i\Delta k \cdot x),
\]

where \( \Delta k = p_1 - p_2, K = (p_1 + p_2)/2 \) and \( \Delta k \cdot x \) stands for the inner product of the four-vectors. The invariant momentum distribution reads as

\[
E \frac{d\eta}{dp} = N_1(p) = \tilde{S}(\Delta k = 0, K = p).
\]

In this Letter we shall utilize the ‘hydrodynamical’ normalization of the Wigner-functions,

\[
\int \frac{dp}{E} d^4x S(x, p) = \int \frac{dp}{E} N_1(p) = \langle n \rangle,
\]

where \( \langle n \rangle \) is the mean multiplicity. The two-particle BECF-s are prescribed as

\[
C_2(p_2, p_2) \approx 1 + \frac{|\tilde{S}(\Delta k, K)|^2}{S(0, p_1)S(0, p_2)}, \approx 1 + \frac{|\tilde{S}(\Delta k, K)|^2}{|S(0, K)|^2},
\]

as was presented e.g. in ref. [20, 7]. The precision of the last approximation was estimated to be of the order of 5 % as discussed in Ref. [34]. In this Letter final state interactions are neglected and a completely chaotic particle emission is assumed.

One can show [35] that the higher order Bose-Einstein correlation functions,

\[
C_n(p_1, p_2, ..., p_n) = \frac{N_n(p_1, p_2, ..., p_n)}{N_1(p_1)N_1(p_2)...N_1(p_n)}
\]

(5)
are given in terms of the Fourier-transformed Wigner-functions as

\[ C_n(p_1, p_2, ..., p_n) = \sum_{\sigma^n} \prod_{i=1}^{n} \tilde{S}(i, \sigma_i) \prod_{i=1}^{n} \tilde{S}(i, i) = \sum_{\sigma^n} \prod_{i=1}^{n} \tilde{S}(i, \sigma_i) = \sum_{\sigma^n} \prod_{i=1}^{n} \tilde{s}(i, \sigma_i), \tag{6} \]

where the summation is over all the \( \sigma^n \) permutations of \( n \) indexes,

\[ \tilde{S}(i, \sigma_i) = \tilde{S}(K_{i,\sigma_i}, \Delta k_{i,\sigma_i}), \quad \text{and} \quad \tilde{s}(i, \sigma_i) = \frac{\tilde{S}(i, \sigma_i)}{\tilde{S}(i, i)}, \tag{7} \]

\[ K_{i,\sigma_i} = \frac{p_i + p_{\sigma_i}}{2}, \quad \text{and} \quad \Delta k_{i,\sigma_i} = p_i - p_{\sigma_i}. \tag{8} \]

Note the distinction between \( \sigma^n \), which stands for the set of the permutations of the numbers 1, 2, ..., \( n \) and the \( \sigma_i \) (subscript \( i \)), which indicates the number replacing the index \( i \) in a given permutation from the set \( \sigma^n \). Note also that \( \tilde{s}(i, \sigma_i) = \tilde{S}(\sigma_i, \sigma_i) \tilde{s}^*(\sigma_i, i) \neq \tilde{s}^*(\sigma_i, i) \), although the relationship \( \tilde{S}(i, \sigma_i) = \tilde{S}^*(\sigma_i, i) \) is satisfied.

In the core/halo model, the following assumptions are made:

**Assumption 0:** The emission function does not have a no-scale, power-law like structure. This possibility was discussed and related to intermittency in ref. [37].

**Assumption 1:** The bosons are emitted either from a central part or from the surrounding halo. Their emission functions are indicated by \( S_c(x, p) \) and \( S_h(x, p) \), respectively. According to this assumption, the complete emission function can be written as

\[ S(x, p) = S_c(x, p) + S_h(x, p), \tag{9} \]

using the hydrodynamic normalization of the Wigner functions.

**Assumption 2:** We assume that the emission function which characterizes the halo changes on a scale \( R_H \) which is larger than \( R_{\text{max}} \approx h/Q_{\text{min}} \), the maximum length-scale resolvable [18] by the intensity interferometry microscope. However, the smaller central part of size \( R_c \) is assumed to be resolvable, \( R_H > R_{\text{max}} > R_c \). This inequality is assumed to be satisfied by all characteristic scales in the halo and in the central part, e.g. in case the side, out or longitudinal components [33, 21] of the correlation function are not identical.

**Assumption 3:** The momentum-dependent core fraction \( f_c(i) = N_c(p_i)/N_1(p_i) \) varies slowly on the relative momentum scale given by the correlator of the core \( \tilde{s}_c(1, 2)\tilde{s}_c(2, 1) \).

According to this smoothness assumption, which was only implicitly made in ref. [15], the relative momentum dependence of the core fraction can be neglected on the relative
momentum scales on which the correlator of the core is changing. If a typical core radius parameter in a given direction is indicated by $R_c$ then the slow variation of $f_c(i)$ implies that $f_c(p_i)f_c(p_j) \simeq f_c^2(K_{ij})$ if $|p_i - p_j| < \hbar/R_c$.

The emission function of the center and that of the halo are normalized here as

$$\int d^4x \frac{d^4p}{E} S_c(x,p) = \langle n \rangle_c, \quad \text{and} \quad \int d^4x \frac{d^4p}{E} S_h(x,p) = \langle n \rangle_h,$$

where the subscripts $c, h$ index the contributions by the central and the halo parts, respectively. In ref [18], the above Assumptions were formulated with the help of the Wigner-functions normalized to 1, and a core-fraction $f_c = \langle n \rangle_c/\langle n \rangle$ was used, while in this paper the momentum-dependent core fraction $f_c(p)$ is utilized. One finds that

$$N_1(p) = N_c(p) + N_h(p), \quad \text{and} \quad \langle n \rangle = \langle n \rangle_c + \langle n \rangle_h. \quad (11)$$

Note, that in principle the core as well as the halo part of the emission function could be decomposed into more detailed contributions. In case of pions and NA44 acceptance [38], one can separate the contribution of various long-lived resonances as

$$S_h(x,p) = \sum_{r=\omega,\eta,\eta',K_0} S^{(r)}_{\text{halo}}(x,p) \quad \text{and} \quad N_h(p) = \sum_{r=\omega,\eta,\eta',K_0} N^{(r)}_{\text{halo}}(p). \quad (12)$$

For our considerations, this separation is indifferent. According to our assumptions, the Fourier - transformed Wigner-functions characterizing the (full) halo part are narrower than the two-particle momentum resolution $Q_{\text{min}}$, for which we typically may take the value of cca. 10 MeV. It is important to keep in mind that the halo part of the emission function is defined with respect to $Q_{\text{min}}$. For example, if $Q_{\text{min}} = 10 - 15$ MeV, the decay products of the $\omega$ resonances can be taken as parts of the halo [38], however, should the experimental resolution and the error bars on the measurable correlation function decrease significantly below $\hbar/\Gamma_\omega = 8$ MeV, the decay products of the $\omega$ resonances would contribute to the core. See refs. [38, 18] for a more complete discussion on this specific topic.

The measured $\tilde{S}_h(i,\sigma_i)$ is thus vanishing if $i \neq \sigma_i$ at the given measured relative momentum $Q_{i,\sigma_i} > Q_{\text{min}}$. Note that $Q_{i,\sigma_i}$ can be one, two- or three-dimensional quantity, e.g. any reasonable measure of $\Delta k_{i,\sigma_i}$. However, it is important to observe that $\tilde{S}_h(i,i)$ contributes to the spectrum, and is not affected by the two-particle momentum resolution:

$$\tilde{S}_h(i,\sigma_i) = \delta_{i,\sigma_i} \tilde{S}(i,i), \quad (13)$$

$$\tilde{s}(i,\sigma_i) = \delta_{i,\sigma_i} + (1 - \delta_{i,\sigma_i}) f_c(i) \tilde{s}_c(i,\sigma_i), \quad (14)$$

$$f_c(i) = f_c(p_i) = N_c(p_i)/N_1(p_i). \quad (15)$$
Thus the Bose-Einstein correlation function $C_n(1, 2, \ldots, n) = C(p_1, p_2, \ldots, p_n)$ reads as

$$C_n(1, 2, \ldots n) = \sum_{\sigma_n} \prod_{i=1}^n [\delta_{c, \sigma_i} + (1 - \delta_{i, \sigma_i}) f_c(i) \tilde{s}_c(i, \sigma_i)]$$

(16)

Let us introduce $\rho^n$ to stand for those permutations of $(1, \ldots, n)$ which are mixing all the numbers from 1 to $n$ and let us indicate by $\rho_i$ the value which is replaced by $i$ in a given permutation belonging to the set of permutations $\rho^n$. (Superscript indexes a set of permutations, subscript stands for a given value). Then we have $\rho_i \neq i$ for all values of $i = 1, \ldots, n$, while $\sigma_i = i$ type of replacements are allowed.

Thus the general expression for $C_n(1, \ldots, n)$ reads as

$$C_n(1, \ldots, n) = 1 + \sum_{j=2}^n \sum_{i_1, \ldots, i_j=1}^{\rho^n} \prod_{k=1}^j f_c(i_k) \tilde{s}_c(i_k, i_{p_k}).$$

(17)

Here $\sum'$ indicates that the summation should be taken over those set of values of the indices which do not contain any value more than once.

Let us explicitly write out the above expression until the $l = 4$ terms:

$$C_n(1, \ldots, n) = 1 + \sum_{i \neq j=1}^n f_c(i) f_c(j) [\tilde{s}_c(i, j) \tilde{s}_c(j, i)] +$$

$$+ \sum_{i \neq j \neq k=1}^n f_c(i) f_c(j) f_c(k) [\tilde{s}_c(i, j) \tilde{s}_c(j, k) \tilde{s}_c(k, i) + \tilde{s}_c(i, k) \tilde{s}_c(k, j) \tilde{s}_c(j, i)] +$$

$$+ \sum_{i \neq j \neq k \neq l=1}^n f_c(i) f_c(j) f_c(k) f_c(l) \times$$

$$\times [\tilde{s}_c(i, j) \tilde{s}_c(j, k) \tilde{s}_c(k, l) \tilde{s}_c(l, i) + \tilde{s}_c(i, j) \tilde{s}_c(j, l) \tilde{s}_c(l, k) \tilde{s}_c(k, i) +$$

$$+ \tilde{s}_c(i, k) \tilde{s}_c(k, j) \tilde{s}_c(j, l) \tilde{s}_c(l, i) + \tilde{s}_c(i, k) \tilde{s}_c(k, l) \tilde{s}_c(l, j) \tilde{s}_c(j, i) +$$

$$+ \tilde{s}_c(i, l) \tilde{s}_c(l, j) \tilde{s}_c(j, k) \tilde{s}_c(k, i) + \tilde{s}_c(i, l) \tilde{s}_c(l, k) \tilde{s}_c(k, j) \tilde{s}_c(j, i) +$$

$$+ \tilde{s}_c(i, j) \tilde{s}_c(j, k) \tilde{s}_c(k, l) \tilde{s}_c(l, i) + \tilde{s}_c(i, j) \tilde{s}_c(j, k) \tilde{s}_c(k, l) \tilde{s}_c(l, j)] + \ldots.$$ 

(18)

There are 44 (and 265) terms in the next part which fully mixes 5 (and 6) different indexes for any given fixed value of $i \neq j \neq k \neq l \neq m(\neq n)$, respectively.

For the two-particle case, we recover the earlier result given in Ref. [18]. Let us proceed very carefully at this point, so that the role of Assumption 3 could be made transparent. If only two particles are present, the above equation reduces to

$$C_2(1, 2) = 1 + f_c(1) f_c(2) \tilde{s}_c(1, 2) \tilde{s}_c(2, 1).$$

(19)
If Assumption 3 is also satisfied by some experimental data set, then the resulting formula becomes particularly simple:

$$C(\Delta k_{12}, K_{12}) = 1 + \lambda_*(K_{12}) \frac{\left| \tilde{S}_c(\Delta k, K) \right|^2}{\tilde{S}_c(0, p_1) \tilde{S}_c(0, p_2)}, \quad \simeq 1 + \lambda_*(K_{12}) \frac{\left| \tilde{S}_c(\Delta k, K) \right|^2}{\tilde{S}_c(0, K_{12})},$$

(20)

where the effective intercept parameter \( \lambda_*(K_{12}) \) is given as

$$\lambda_*(K_{12}) = \left[ \frac{N_c(K_{12})}{N_1(K_{12})} \right]^2.$$  

As emphasized in Ref. [18], this effective intercept parameter (\( \neq \) exact intercept parameter at \( Q = 0 \) MeV) shall in general depend on the mean momentum of the observed boson pair, which within the errors of \( Q_{\text{min}} \) coincides with any of the on-shell four-momentum \( p_1 \) or \( p_2 \). Thus one obtains the core/halo interpretation of the two-particle correlation function: the intercept parameter \( \lambda_*(K) \) measures the momentum dependent core fraction and the relative momentum dependent part of the two-particle correlation function is determined by the core. Note, that in high energy heavy ion collisions the momentum dependence of the \( \lambda_*(K) \) parameter is very weak, actually, within the errors \( \lambda_*(K) \) is constant for the NA44 data analyzed in ref. [18]. However, the validity of this Assumption 3 has to be checked experimentally for each data set, by determining the momentum dependence of the \( \lambda(K) \) parameter of the two-particle correlation function.

Within the core/halo picture, the \( n \)-particle correlation function has a simple form if all the \( n \) momenta are approximately equal, i.e. \( |\Delta k_{i, \sigma_i}| \leq Q_{\text{min}} \) for all \( i \neq \sigma_i \), and this situation is denoted by \( p_1 \simeq p_2 \simeq \ldots \simeq p_n \simeq P \). One obtains that

$$C_n(p_1 \simeq p_2 \simeq \ldots \simeq p_n \simeq P) = 1 + \sum_{j=1}^{n} f_c(P)^j \binom{n}{j} \alpha_j$$

where \( \alpha_j \) stands for the number of fully mixing permutations of \( j \) indexes, i.e. the number of permutations in \( \rho^j \). The first few values of \( \alpha_j \) are given as

$$\alpha_1 = 0, \quad \alpha_2 = 1, \quad \alpha_3 = 2, \quad \alpha_4 = 9, \quad \alpha_5 = 44, \quad \alpha_6 = 265.$$  

(22)

These values can be obtained from a recurrence relation, as follows. Let us indicate the number of permutations that completely mix exactly \( j \) non-identical elements by \( \alpha_j \). There are exactly \( \binom{n}{j} \) different ways to choose \( j \) different elements from among \( n \) different elements. Since all the \( n! \) permutations can be written as a sum over the fully mixing permutations, the counting rule yields \( n! = 1 + \sum_{j=1}^{n} \binom{n}{j} \alpha_j \), which can be rewritten as a recurrence relation for \( \alpha_j \):

$$\alpha_n = n! - 1 - \sum_{j=1}^{n-1} \binom{n}{j} \alpha_j.$$  

(23)
We have the following explicit expressions for the first few intercept parameters:

\[
\lambda^{*}_{2}(P) = f_{c}(P)^2 \\
\lambda^{*}_{3}(P) = 3f_{c}(P)^2 + 2f_{c}(P)^3 \\
\lambda^{*}_{4}(P) = 6f_{c}(P)^2 + 8f_{c}(P)^3 + 9f_{c}(P)^4 \\
\lambda^{*}_{5}(P) = 10f_{c}(P)^2 + 20f_{c}(P)^3 + 45f_{c}(P)^4 + 44f_{c}(P)^5 \\
\lambda^{*}_{6}(P) = 15f_{c}(P)^2 + 40f_{c}(P)^3 + 135f_{c}(P)^4 + 264f_{c}(P)^5 + 265f_{c}(P)^6
\]

In general, the intercept parameter of the \(n\)-particle correlation function reads as

\[
\lambda^{*}_{n}(P) = \sum_{j=1}^{n} f_{c}(P)^j \left( \begin{array}{c} n \\ j \end{array} \right) \alpha_j
\]

where \(\alpha_j\) is defined by the recurrence given in Eq. (23). These expressions then relate the effective intercept parameter of the \(n\)-particle correlation function to the effective intercept parameter of the two-particle correlation function, and could be thus checked experimentally. For that type of test, a few warnings should be made: The intercept parameters \(\lambda^{*}_{n}(P)\) explicitly depend on the momentum of the particles in the region where all the \(n\) particles have approximately the same momentum. Thus averaging over the transverse mass and the rapidity of the particles to improve statistics is in principle not allowed. Further, the effective intercept parameter \(\lambda^{*}_{n}(P)\) differs from the exact analytical value of the \(n\)-particle correlation function which is \(n!\) in our picture, due to the fact that we have neglected possible partial coherence of the particle emitting source. Our motivation for ignoring possible partial coherence was to find a relatively simple limiting case, and also we are aware of the fact that the pure quantum statistical relationship between the second order and the higher order correlation functions has been shown not to be consistent with the available UA1 data [12]. In case of heavy ion reactions, resonance halo seems to be describing the drop of the effective intercept parameter from the value of 2 both in numerical simulations [22] and in analytical approaches [18, 13]. Finally, we stress that the values for the effective intercept parameters \(\lambda^{*}_{n}(P)\) may depend on the two-particle momentum resolution \(Q_{\text{min}}\), on the error distribution on the correlation function, as well as on the Gaussian or non-Gaussian structure for the Fourier-transformed Wigner-function of the core, \(\tilde{S}_{c}(K, \Delta k)\).

Note also that the general result for the correlation function in the core/halo model coincides with a particular limiting case of the expressions obtained by Cramer and Kadija for the correlation functions of the order 1, 2, 3, 4 and 5, namely, the case when particle mis-identification (contamination) is taken into account but the source is assumed to have
no partially coherent component. Thus our results can be applied also to the evaluation of the intercept parameter of higher order correlation functions for data contaminated by unidentified or erroneously identified particles.

**Prediction of higher order intercept parameters for NA44 data.** Equations (24-28) can also be used to predict the intercept parameter of higher order correlation functions from the measured intercept parameter of the second order correlation function. The published NA44 correlation data indicate that $\lambda_{2,*} \simeq 0.52 \pm 0.02$ approximately independently of the transverse mass for both the low $p_t$ and the high $p_t$ data sample for pion pairs in a central $S + Pb$ reaction at 200 A GeV. If we interpret this $\lambda_{2,*}$ in the core/halo model, we obtain that $f_c = 0.72$ is approximately independent of the transverse mass in the NA44 acceptance range. Thus, the core/halo model predicts the following measurable intercept parameters for the higher order correlation functions of identified pions in 200 AGeV $S + Pb$ reactions in the NA44 acceptance:

$$
\lambda_{*,3}(P) = 2.3 \quad \lambda_{*,4}(P) = 8.6 \quad \lambda_{*,5}(P) = 33.4 \quad \lambda_{*,6}(P) = 148.0 \quad (30)
$$

It should be emphasized that the NA44 data constitute a good sample to test the core/halo model on higher order correlation functions, given the experimentally observed momentum independence of the $\lambda_{2,*}$ parameter [31].

**In summary,** we have evaluated the higher order correlation functions in the core/halo model analytically. A closed form of the correlation function of arbitrary high order is given by eq. (17) in terms of the momentum-dependent core fractions and in terms of permutations that completely mix a set of indexes such that $\rho_i \neq i$. A recurrence relation has been found which allows for the evaluation of the measured intercept parameter of the $n$-particle correlation function for arbitrary high orders in a very efficient manner. We emphasized that the separation of the core from the halo is dependent on the relative momentum resolution of the experiment and pointed out a formal analogy between the correlation functions of the core/halo model and those of completely chaotic sources with particle mis-identification. Our results allow for a prediction of the higher order correlation functions if the basic building block, the amplitude $\tilde{s}(i, j)$ is determined experimentally as a function of the mean and the relative momentum of the particle pair. If this quantity is real, it can be determined from a detailed analysis of the two-particle correlation function. If the imaginary part of $\tilde{s}(i, j)$ is not negligible, a simultaneous analysis of second and third order correlation functions is necessary to extract the building block of higher order correlations.

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