Physics Letters B 762 (2016) 57–63

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Multipion correlations induced by isospin conservation of coherent emission

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ARTICLE INFO

Article history:
Received 1 June 2016
Received in revised form 3 September 2016
Accepted 13 September 2016
Available online 15 September 2016
Editor: J.-P. Blaizot

ABSTRACT

Recent measurements have revealed a significant suppression of multipion Bose–Einstein correlations in heavy-ion collisions at the LHC. The suppression may be explained by postulating coherent pion emission. Typically, the suppression of Bose–Einstein correlations due to coherence is taken into account with the coherent state formalism in quantum optics. However, since charged pion correlations are most often measured, the additional constraint of isospin conservation, which is absent in quantum optics, needs to be taken into account. As a consequence, correlations emerge between pions of opposite charge. A calculation of the correlations induced by isospin conservation of coherent emission is made for two, three- and four-pion correlation functions and compared to the data from the LHC.
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1. Introduction

Data from high energy collisions present a unique opportunity to study the possibility of quantum coherence at very high temperatures. Bose–Einstein (BE) correlations of identical bosons are often used to search for coherent emission. In the presence of coherent emission, the strength of BE correlations is expected to be suppressed. A 6σ suppression of four-pion BE correlations was recently found in Pb–Pb collisions at the LHC with ALICE [1].

Multibody Coulomb correlations have also been proposed as a potential source of the suppression. However, such correlations are unlikely to decrease with K_T, in contrast with the observed decrease of the suppression at high K_T. An important additional feature of BE correlations in high multiplicity collisions is their robustness to background correlations unrelated to quantum statistics (QS) and final-state interactions (FSI). The QS and FSI correlations for large emission volumes occur in a very narrow interval in relative momentum while all known background correlations produce easily distinguished broad correlations.

While two-pion BE correlations are most often measured experimentally, they are insufficient to search for coherent emission. The unknown freeze-out distribution of particles produced in collisions causes two noteworthy uncertainties in a two-pion analysis. The first being the fraction of pions from short- compared to long-lived emitters, which characterizes a dilution of BE correlations. Secondly, the computation of FSI is done with an assumed freeze-out distribution. Both of these uncertainties make it practically impossible to determine the presence of coherent emission from two-pion correlations alone. However, both uncertainties were found to largely cancel in the comparison of measured to *built* multipion BE correlations [1,2].

The effect of pion coherence on BE correlations is typically incorporated using the coherent state formalism of quantum optics [3]. However, the fact that *charged* pion correlations are measured necessitates an extension of the quantum optics approach using the super-selection rule [4–8]. The resulting correlations induced by isospin conservation of the coherent component occurs between all pion species. Isospin conservation of the chaotic component also induces additional correlations in scenarios where the emission duration of the source is short [9]. In this letter, expressions are derived for three- and four-pion correlations stemming from isospin conservation of pion coherent states. Calculations are presented for four different mixed-charge correlation functions and compared to the LHC data.

2. Formalism

The formalism of pion coherent states obeying the super-selection rule and their application to BE correlations is given in detail in Refs. [7] and [8]. Coherent states are taken in the standard form as given by Eq. 8 of Ref. [7] and is also the same as various investigations of disoriented chiral condensates [10–12]. It is assumed that particle production can be split into a chaotic and a single coherent component, whose annihilation operators at a given momentum are given by $b(p)$ and $d(p)$, respectively.

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http://dx.doi.org/10.1016/j.physletb.2016.09.017
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$$a_i(p) = b_i(p) + e_i d(p) \tag{1}$$

$$e_0 = \cos(\theta) \tag{2}$$

$$e_\pm = \frac{\sin\theta}{\sqrt{2}} e^{\pm\phi}. \tag{3}$$

The subscript i denotes the pion type (π^+, π^-, π^0) for which one commonly introduces a unit vector, $\mathbf{e}$, in three-dimensional isospin space. It is noteworthy to point out that the coherent component of pion production is assumed to conserve isospin independently from the chaotic component. The resulting isospin conservation induced correlations are to be understood as a theoretical maximum which can be diminished if the two components do not independently conserve isospin. The single particle inclusive momentum densities are given in the usual way

$$\begin{aligned} N_i^{(1)}(p) &\equiv E_{\mathbf{p}} \frac{d^3 N_i}{d^3\mathbf{p}} = \langle a_i^\dagger(p) a_i(p)\rangle \\ &= \langle b_i^\dagger(p) b_i(p)\rangle + \langle |e_i|^2\rangle \langle d^\dagger(p) d(p)\rangle \\ &= N_{\rm ch}^{(1)}(p) + N_{\rm coh}^{(1)}(p), \end{aligned} \tag{4}$$

where $E_{\mathbf{p}} = \sqrt{m^2 + \mathbf{p}^2}$. The chaotic and coherent components are given by ch and coh, respectively. An averaging over all possible orientations of the isospin vector is done in order to compute all final observables. The following integrals over the isospin vector will be needed to evaluate the two-, three-, and four-pion correlation functions.

$$\langle |e_+|^n\rangle = \frac{1}{4\pi} \int d\cos(\theta) d\phi \left[\frac{\sin\theta}{\sqrt{2}}\right]^n \tag{5}$$

$$\langle |e_+|^2\rangle = \langle |e_-|^2\rangle = \langle |e_0|^2\rangle = \frac{1}{3} \tag{6}$$

$$\langle |e_+|^4\rangle = \frac{2}{15} \quad \langle |e_+|^6\rangle = \frac{2}{35} \quad \langle |e_+|^8\rangle = \frac{8}{315} \tag{7}$$

Integrals which contain a mixture of e_+ and e_- are identical to the ones given above. The total number of pions which are radiated from the classical source at momentum $\mathbf{p}$ is given by $|d(\mathbf{p})|^2$ while the coherent fraction of pions is defined as $G(p) \equiv \frac{N_{\rm coh}^{(1)}(p)}{N^{(1)}(p)} = \frac{1}{3}|d(\mathbf{p})|^2$.

It is convenient to introduce the single particle Wigner function, split into chaotic and coherent components

$$f_{\mathbf{e},i}(x,p) = f_{\rm ch}(x,p) + |e_i|^2 f_{\rm coh}(x,p), \tag{8}$$

which provide the following two important links between the expectation values of the pionic field operators and an integration over the freeze-out hypersurface ($\sigma_{\rm out}$),

$$\begin{aligned} \langle b_i^\dagger(p_1) b_i(p_2)\rangle &= \int_{\sigma_{\rm out}} d\sigma_\mu p^\mu f_{\rm ch}(x,p) e^{-iqx} \\ &\equiv T_{12} e^{-i\Phi_{12}} \sqrt{[1-G(p_1)][1-G(p_2)] N_i^{(1)}(p_1) N_i^{(1)}(p_2)}, \end{aligned} \tag{9}$$

$$\begin{aligned} \langle d_i^\dagger(p_1) d_i(p_2)\rangle &= \langle |e_i|^2\rangle \int_{\sigma_{\rm out}} d\sigma_\mu p^\mu f_{\rm coh}(x,p) e^{-iqx} \\ &\equiv t_{12} e^{-i\phi_{12}} \sqrt{G(p_1) G(p_2) N_i^{(1)}(p_1) N_i^{(1)}(p_2)}, \end{aligned} \tag{10}$$

where $q = p_1 - p_2$ and $p = (p_1 + p_2)/2$. The pair exchange magnitudes of the chaotic and coherent components are denoted by T_{ij} and t_{ij}, respectively. For the expectation value of two or more coherent pions with an imbalance of operators at momentum p_1 and p_2, we have the relation

$$\begin{aligned} &\langle d_{\omega_1}^\dagger(p_1) d_{\omega_1}(p_2) \cdot d_{\omega_2}^\dagger(p_3) d_{\omega_2}(p_3) \cdot \ldots \cdot d_{\omega_n}^\dagger(p_n) d_{\omega_n}(p_n)\rangle \\ &= \langle \prod_\gamma^n |e_{\omega_\gamma}|^2\rangle \left[\int d\sigma_\mu p^\mu f_{\rm coh}(x,p) e^{-iq_{12}x}\right] \langle d^\dagger(p) d(p)\rangle^{n-1} \end{aligned} \tag{11}$$

3. Three- and four-pion quantum statistics correlation functions

The multipion inclusive momentum density distributions is given in the usual way as

$$\begin{aligned} N_{\omega_1\ldots\omega_n}^{(n)}(p_1,\ldots,p_n) &= \left[\prod_{\alpha=1}^n E_{\mathbf{p}_\alpha}\right] \frac{d^{3n} N_{\omega_1\ldots\omega_n}}{\prod_{\alpha=1}^n d^3\mathbf{p}_\alpha} \\ &= \langle \prod_{\alpha=1}^n a_{\omega_\alpha}^\dagger(p_\alpha) a_{\omega_\alpha}(p_\alpha)\rangle, \end{aligned} \tag{12}$$

where ω represents the set of n elements, each of which denote a particular type of pion. For example, the set ω in the case of the $\pi^+\pi^+\pi^-$ distribution is given by $\omega_1 = \pi^+$, $\omega_2 = \pi^+$, $\omega_3 = \pi^-$.

Experimentally, the multipion distributions are often projected onto the Lorentz invariant relative momentum and average pair transverse momentum defined by

$$Q_n = \sqrt{-\sum_{i=1}^{n-1} \sum_{j=i+1}^n (p_i - p_j)^2}, \tag{13}$$

$$K_{{\rm T}n} = |\sum_{i=1}^n \mathbf{p}_{{\rm T},i}|/n. \tag{14}$$

The three- and four-pion QS distributions are decomposed into several components,

$$N_{ijk}^{(3)} \equiv I_1 + I_2 + I_3, \tag{15}$$

$$N_{ijkl}^{(4)} \equiv J_1 + J_2 + J_3 + J_4, \tag{16}$$

where the I_1 and J_1 will be defined to contain the conventional QS correlations as prescribed by quantum optics. The other components arise from the constraint of isospin conservation of coherent emission. Neglecting FSI, the components for three-pion correlations are given in Eqs. (17)–(19) while those for four-pion correlations are given in the appendix.

$$\begin{aligned} I_1 &= N_i^{(1)}(p_1) N_j^{(1)}(p_2) N_k^{(1)}(p_3) \\ &+ \sum_\omega \delta_{\omega_\alpha\omega_\beta} \Big[|\langle b_{\omega_\alpha}^\dagger(p_\alpha) b_{\omega_\alpha}(p_\beta)\rangle|^2 \\ &+ 2\Re\Big\{\langle b_{\omega_\alpha}^\dagger(p_\alpha) b_{\omega_\alpha}(p_\beta)\rangle \langle d_{\omega_\alpha}^\dagger(p_\beta) d_{\omega_\alpha}(p_\alpha)\rangle\Big\}\Big] N_{\omega_\gamma}^{(1)}(p_\gamma) \\ &+ 2\delta_{ijk}\Big[\Re\Big\{\langle b_i^\dagger(p_1) b_i(p_2)\rangle \langle b_i^\dagger(p_2) b_i(p_3)\rangle \langle b_i^\dagger(p_3) b_i(p_1)\rangle\Big\} \\ &+ 3\Re\Big\{\langle d_i^\dagger(p_1) d_i(p_2)\rangle \langle b_i^\dagger(p_2) b_i(p_3)\rangle \langle b_i^\dagger(p_3) b_i(p_1)\rangle\Big\}\Big], \end{aligned} \tag{17}$$

$$\begin{aligned} I_2 &= \sum_\alpha \langle b_{\omega_\alpha}^\dagger(p_\alpha) b_{\omega_\alpha}(p_\alpha)\rangle \\ &\times \Big[\langle \prod_{\epsilon\in\omega\backslash\{\alpha\}} d_{\omega_\epsilon}^\dagger(p_\epsilon) d_{\omega_\epsilon}(p_\epsilon)\rangle - \prod_{\epsilon\in\omega\backslash\{\alpha\}} \langle d_{\omega_\epsilon}^\dagger(p_\epsilon) d_{\omega_\epsilon}(p_\epsilon)\rangle\Big] \\ &+ 2\sum_\omega \delta_{\omega_\alpha\omega_\beta} \Re\Big\{\Big[\langle d_{\omega_\alpha}^\dagger(p_\alpha) d_{\omega_\gamma}^\dagger(p_\gamma) d_{\omega_\alpha}(p_\beta) d_{\omega_\gamma}(p_\gamma)\rangle \\ &- \langle d_{\omega_\alpha}^\dagger(p_\alpha) d_{\omega_\alpha}(p_\beta)\rangle \langle d_{\omega_\gamma}^\dagger(p_\gamma) d_{\omega_\gamma}(p_\gamma)\rangle\Big] \langle b_{\omega_\alpha}^\dagger(p_\beta) b_{\omega_\alpha}(p_\alpha)\rangle\Big\}, \end{aligned} \tag{18}$$

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