PROBLEMS OF PARTICLE INTERFEROMETRY*

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The problems of particle interferometry arising from the finite size of the particle sources, various types of the sources, resonances, final state interactions, rescattering and absorption effects are discussed.

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1. Introduction

The particle interferometry is understood as the measurement of the space-time characteristics of the production process with the help of narrow particle correlations. Thus in astronomy the two-photon interferometry (HBT effect [1]) is used to measure angular radii of stars. In this case the interference correlations appear solely due to the effect of quantum statistics (QS), *i.e.* due to the identity of photons. The interferometry in particle physics is more complicated due to the effect of final state interactions (FSI). On the other hand, this effect allows one to measure the space-time characteristics of the production process even with the help of nonidentical particles. Both the effects of QS and FSI have been widely used to measure space-time dimensions of the production region of various particles and nuclear fragments in multiparticle processes (see, *e.g.*, reviews [2–4]).

It should be noted that there is a principal difference between interferometry in astronomy and in particle physics. In astronomy the dependence of the coincidence rate on the distance between the photon detectors is measured (space-time variant) while in particle physics the momentum-energy variant is realized. The interference QS effect in particle physics was first observed in $\bar{p}p$ annihilations as an enhanced production of the pairs of identical

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charged pions with small opening angles (GGLP effect [5]). The interpretation problems prevented however this method to be widely used as a tool for particle interferometry. The situation changed in early seventieth when Kopylov and Podgoretsky [2,6,7] observed the analogy between interferometry in astronomy and particle physics, suggested to study the interference effect in terms of the correlation function (similar to astronomy, where the coincidence rate is normalized with the help of single detector counts to get rid of the unknown gain of the apparatus) and presented a simple space-time parametrization of the correlation function. The K–P method is now commonly used in particle interferometry. The simplicity of this approach is however partly lost due to the delicate problems with the reference distribution required to construct the correlation function, especially in the processes characterized by relatively small space-time parameters (*e.g.*, in e^+e^- annihilations at LEP [8–10]).

Below we consider some other problems of particle interferometry related to the finite size of the particle sources, various types of the sources, resonances, FSI, rescattering and absorption effects.

2. Finite-size one-particle sources

In the original K-P model the noninteracting identical particles were assumed to be emitted independently by heavy one-particle point-like sources distributed in some space-time region. For the correlation function of two identical pions, defined as the ratio of the differential production cross section to the one which would be observed in the case of no interference effects, this model yields the well-known result [2]:

$$R(p_1,p_2) = 1 + \int W(x_A,x_B)\cos(qx)\, d^4x_A\, d^4x_B \equiv 1 + \langle\cos(qx)\rangle\,. \qquad (1)$$

Here $q \equiv \{q_0,\vec{q}\} = p_1 - p_2$ and $x \equiv \{t,\vec{x}\} = x_A - x_B$ are the differences of the pion 4-momenta and the 4-coordinates of the emission points and $W(x_A,x_B) = w(x_A)w(x_B)$ is the normalized distribution of these points. Note that in the case of identical particles with nonzero spin the sign of the interference term is positive (negative) for even (odd) total spin S of the pair.

The characteristic feature of the correlation function (1) is the presence of the interference maximum at small $|q|$ changing to a horizontal plateau at sufficiently large $|q|$. *E.g.*, assuming that

$$w(x_I;p) \propto \exp\left(-\frac{\vec{x}_I^{\,2}}{2r_0^2} - \frac{t_I^2}{2\tau_0^2}\right) \qquad (2)$$

effectively describes the space and time limitation of the production process, we get $R(p_1, p_2) = 1 + \exp(-r_0^2\vec{q}^2 - \tau_0^2 q_0^2)$. It is quite possible that the particles are emitted by two or more types of the sources characterized by different space-time parameters [11]; a special case is the admixture of multiparticle sources [12] (emitting particles in one and the same quantum state and thus not interfering with each other; similar effect results from the so called "coherent states" [13]). In the presence of very "wide" or very "narrow" sources, the parametrization of the correlation function with one set of the space-time parameters would lead to an effective suppression of the interference effect. This can be taken into account by an additional parameter λ (sometimes called "degree of incoherence") multiplying the interference term.

If three or more identical pions are produced, the symmetrization over all pion 4-momenta leads to a large number of the interference terms [2, 12, 14, 15]. *E.g.*, the correlation function of three identical pions emitted by the sources of the same type contains four interference terms, the three of them reducing to the usual pair correlations, while the fourth one is of specific three-particle origin (it is substantial only if all pions have nearby momenta). It is easy to see that at sufficiently high energies, when the typical distance between pions in momentum space is much larger than the width of the interference region, the difference between the two-pion correlation functions in the multi-pion and two-pion events is small and vanishes with the increasing mean energy of relative motion of the pions. In such situation the simple two-particle K–P approach is justified.

Consider now the effect of a finite size of the source. Assume that a heavy source excited at a space-time point $x_A^{(0)}$ with lifetime τ and velocity $\vec{\beta}$ emits a pion with 4-momentum $p \equiv \{\omega, \vec{p}\}$ at the point $x_A = \{t_A, \vec{x}_A^{(0)} + \vec{\beta}(t_A - t_A^{(0)})\}$ according to the emission amplitude $T(x_A^{(0)}; \vec{\beta}) \exp\left(-\frac{t_A - t_A^{(0)}}{2\tau\gamma}\right)$ $\times \theta(t_A - t_A^{(0)}) u(p; \vec{\beta}) e^{ipx_A}$, where γ is the Lorentz factor of the source. The function $u(p; \vec{\beta}) = \bar{u}[p^*(p; \vec{\beta})]$, taking into account a finite size of the source, depends on the velocity $\vec{\beta}$ through the Lorentz transformation $p \rightarrow p^*$ to the source rest frame. Usually the characteristic space and time distances between the emission points are much larger than the characteristic coherence length $L \sim 1/\langle|\vec{p}|\rangle$ and time $T \sim 1/\langle\omega\rangle$ (see, however, [16]),

$$\langle|\vec{x}\,|\rangle \gg L\,, \qquad \langle|x_0|\rangle \gg T\,. \tag{3}$$

It is easy to show that, in this case, the correlation function can be again written in the form (1), but now the distribution of the emission points generally depends on the pion pair 4-momentum $2p = p_1 + p_2$: $W(x_A, x_B; p) =$

$w(x_A;p)w(x_B;p)$,

$$w(x_A;p) \propto \int |u(p;\vec{\beta})|^2 \exp(-\frac{t_A - t_A^{(0)}}{\tau\gamma})\theta(t_A - t_A^0) \times \left|T\left[t_A^{(0)}, \vec{x}_A - \vec{\beta}(t_A - t_A^{(0)}); \vec{\beta}\right]\right|^2 dt_A^0\, d^3\vec{\beta}. \quad (4)$$

Though obtained in the model of heavy one-particle sources, the result (1), with the momentum-dependent distribution of the emission points (generally not factorized), is of general validity provided that condition (3) is fulfilled. *E.g.*, this result can be proved with the help of general formalisms of space-time density matrix [17] or Wigner functions [18].

It follows from Eq. (4) that the correlation between 4-coordinates of the emission points and particle 4-momenta arises from the averaging of the one-particle spectrum over the source velocities. Thus in the case of substantial relative motion of the sources and limited emission momenta (due to finite size of the sources or kinematical constraints) the interfering particles with nearby velocities are mainly produced at nearby space-time points (sources-resonances [6, 19–22], hydrodynamical expansion [18, 23–26], colour string [27, 28]). In such a case the correlation function depends mainly on the invariant variable $Q^2 \equiv -q^2 = \vec{q}^{\,2} - q_0^2 = \vec{q}^{\,*2}$.

3. Sources–resonances

Since pions in high energy multiparticle processes are mainly produced through decays of light resonances, the consideration of their role is especially important for pion interferometry. Consider the model "resonance + particle" in which one of the two identical pions with 4-momenta p_1 and p_2 is produced together with a particle 3 (or a group of particles) in the resonance decay $R \to 1 + 3$ and the other pion and the resonance are emitted at points x_B and x_A by finite size moving sources. On condition (3) and at sufficiently large energy release in the decay: $M - m_1 - m_3 \gg \Gamma/2$ (M and Γ are the resonance mass and width), we get the following expression for the correlation function at small values of $Q < 1/l^*$ [6, 20–22] $R(p_1,p_2) \approx \langle 1 + Re[e^{iqx}/(1 - iy)]\rangle$, where $y = (kq)/(M\Gamma) \equiv (lq) = -\vec{l}^{\,*}\vec{Q}$, $k = p + p_3$ ($l^* \approx p_D/(m_\pi\Gamma)$ is the resonance decay length in the c.m.s. of identical pions, p_D is the decay momentum; $l^* = 3.3 fm$ for ρ). This approximation corresponds to the treatment of the resonance as a classical heavy source with the proper lifetime $1/\Gamma$, often used in Monte Carlo simulations (see, *e.g.*, [29–31]). Note that it overestimates the tail of the correlation function (by $\sim$ 15% in the case of "$\rho + \pi$" model at $x = 0$ [22]). Using this approximation and averaging over the spatial coordinates of the emission points in the c.m.s. of pion pair according to the gaussian with dispersion r_0^2 and over the angle between the vectors $\vec{l}^*$ and $\vec{Q}$, we get [20–22]

$$R(p_1, p_2) \approx 1 + e^{-r_0^2 Q^2} \arctan(l^* Q)/(l^* Q). \tag{5}$$

We see that the finite resonance decay length l^* enlarges the radius of the production region. Fortunately, it substantially influences the slope of the correlation function only at small $Q < \sqrt{3}/l^*$. In fact, it is this circumstance which allows one to determine the size of the "direct" particle production region, despite the extremely small fraction of the pairs of "directly" produced pions. On the other hand, the fit of such correlation function by a single gaussian: $1 + \lambda e^{-r_0^2 Q^2}$ would overestimate the value of r_0 and lead to the suppression factor $\lambda < 1$. Thus, such a fit of the two-pion correlation data obtained at ISR yields $\lambda = 0.45 \pm 0.02$ and $r_0 = 0.88 \pm 0.04$fm with rather bad value of χ^2/NDF ~ 2 [22,32]. The description of this data is substantially improved (χ^2/NDF ~ 1.3) in the case of a two-gaussian fit (the smaller of the fitted radii can be identified with r_0), or — when a superposition of the formulae of the type (5) is used; in both cases the same r_0-value of 0.55 ± 0.08fm is obtained. The fraction of the "direct" pions is estimated to be 0.17 ± 0.09 [22].

Sometimes the effect of two slopes is effectively parametrized by an exponential: $R(p_1, p_2) = 1 + \lambda e^{-rQ}$ or by a more refined parametrization, taking into account the contribution of a multiparticle source ("coherent states"). Thus the recent UA1 data [33] on the correlations of two, three, four and five identical charged pions have been successfully described with approximately the same parameters r and λ in the model with an admixture of the "coherent states". This result was interpreted in Ref. [34] as a strong evidence in favour of the existence of such states. One can get convinced however that the simple model, introducing only an admixture of very "narrow" or very "wide" sources, also well describes this data on correlations of various order with approximately universal parameters r and λ. In fact, the two models substantially differ only in the small-Q region, where the correlation functions are not measured with sufficient accuracy. The UA1 data on multipion correlations thus confirms the selfconsistency of the interferometry approach (found also in other studies) but not more.

4. FSI, rescattering and absorption

For pions the correlation function is mainly determined by the effect of QS, while for nucleons and nuclear fragments the effect of FSI usually dominates. First it was calculated numerically by Koonin [35] for the case of two nonrelativistic protons, while analytical formulae, valid for arbitrary particles produced at sufficiently large distances compared with the range of their strong interaction, were obtained in Ref. [36] (see also [4, 17, 37]). Usually it is assumed that the energy of the relative motion of the multiparticle system is sufficiently large so that the correlation of two particles with a small relative velocity is influenced by their mutual FSI only. The FSI correction ΔT to the nonsymmetrized production amplitude of two noninteracting particles $T_0 \propto u_A(p_1;\vec{\beta}_A)u_B(p_2;\vec{\beta}_B)e^{ip_1x_A+ip_2x_B}$ is determined by the integral over the 4-momenta of the propagating virtual particles. If the momentum dependence of the functions u_A and u_B is sufficiently smooth so that they can be taken out of the integral, the role of FSI is reduced to the substitution of the plane wave in the amplitude T_0 by the Bethe–Salpeter amplitude in the continuous spectrum [17, 37, 38]. At equal emission times $t^* = t^*_A - t^*_B = 0$ in the c.m.s. of the two particles this amplitude coincides up to a phase factor with a stationary solution of the scattering problem $\psi^{(+)}_{-\vec{k}^*}(\vec{x}^*)$ $(\vec{k}^* = \vec{Q}/2)$ having at large $|\vec{x}^*|$ the asymptotics of superposition of the plane and diverging spherical waves. It can be shown that the Bethe-Salpeter amplitude can be substituted by this solution on condition [36] $\tau_0 \ll m\gamma\rho r_0$, where the parameters r_0, τ_0 characterize the distribution of the production points according to Eq. (2), $\rho^2 = r_0^2 + (v\tau_0)^2$, v and γ are the pair velocity and Lorentz factor, m is the reduced mass. This condition is usually satisfied for heavy particles. But even for pions, the use of the equal-time approximation at $r_0 = c\tau_0 = 1fm$ and $\gamma \approx 1$ overestimates the FSI contribution by less than 20% [36]. If further the characteristic distance between the particles in their c.m.s. is larger than the radius of their strong interaction ($\sim 1/m_\rho$ for pions), the solution $\psi^{(+)}$ can be substituted by its asymptotics and the FSI contribution ΔR to the correlation function can be calculated analytically. Neglecting the effect of the Coulomb interaction, we get, *e.g.*, for two pions at $k^* = 0$ [36]:

$$\Delta R(p,p) = \frac{1+\eta}{2\sqrt{\pi}\gamma\rho}\left\{\left|\frac{f(0)}{r_0}\right|^2\left[\sqrt{\pi}A_1 r_0 - \frac{d}{dk^{*2}}\mathrm{Re}\frac{1}{f(k^*)}\right] + 4A_2\mathrm{Re}f(0)\right\}, \tag{6}$$

where $A_1 = \frac{1}{u}\arcsin u$, $A_2 = \frac{1}{2u}\ln\frac{1+u}{1-u}$, $u = \frac{v}{\rho}\sqrt{r_0^2+\tau_0^2}$, $f(k^*)$ is the nonsymmetrized s-wave scattering amplitude, $\eta = 1\,(0)$ for identical (nonidentical) pions. The derivative term represents the first order correction to the spherical wave approximation. For $\pi^+\pi^-$ and $\pi^0\pi^0$ systems a small contribution from the transition $\pi^+\pi^- \longleftrightarrow \pi^0\pi^0$ should be added to ΔR in Eq. (6). In the case of particles with nonzero spin, the contributions corresponding to various values of the total spin should be summed up with the respective weights. We see that ΔR decreases with increasing the particle velocities and the space-time dimensions of the production region. For pions $|f| \ll r$, so that ΔR decreases with the volume $V \sim r^3$ as $V^{-1/3}$. For nucleons, as usual $|f| \gg r$ and $\Delta R \sim V^{-2/3}$. E.g., at $v = 0.7c$, $r_0 = c\tau_0 = 1.5fm$, we have $\Delta R(Q = 0) = -0.07$, 0.13, 0.16 and 15 for the systems $\pi^{\pm}\pi^{\pm}, \pi^0\pi^0, \pi^+\pi^-$ and nn, respectively. Thus the FSI effect on the correlations of identical charged pions is practically negligible while it dominates in nucleon correlations, unless the effective distance between the nucleon production points in their c.m.s. is much larger than their s-wave scattering lengths.

In the case when one of the interfering pions comes from the decay of a resonance, the production amplitude T_0 contains the resonance propagator which cannot be taken out of the integral in ΔT, so that the effect of FSI no more factorizes in the Bethe–Salpeter amplitude. The FSI correction ΔT in the model "resonance + particle" was recently calculated in Ref. [17]. It was shown that in the case of a small space-time distance between the resonance and particle production points, this correction reduces to the well-known logarithmic singularity of the corresponding triangular diagram [39, 40]. This singularity enhances the FSI effect so that, despite the large effective parameter $\gamma v \tau_0 = l^*$ (3.3fm for ρ), ΔR at $Q = 0$ is rather large (~ 0.3 for $\pi^+\pi^-$ in the "$\rho + \pi$" model).

Concerning the role of the Coulomb interaction, it is important only in the region of small momenta $k^* < 2\pi/a_c$ ($\sim 3\,(22)$ MeV/c for two charged pions (two protons); the Bohr radius a_c is $\sim 388\,(58)$fm). If the characteristic distance $\langle|\vec{x}^*|\rangle$ between two particles in their c.m.s. is much smaller than their Bohr radius, the effect of the Coulomb interaction practically factorizes in the function $A_c(k^*)$, representing the modulus squared of the Coulomb wave function at $\vec{x}^* = 0$, *i.e.* $R(p_1,p_2) = A_c(k^*)\tilde{R}(p_1,p_2)$, where $\tilde{R}$ corresponds to to the correlation function of the "neutral" particles. At $k^* \to 0$ the Coulomb repulsion (attraction) forces the correlation function to $0\,(\infty)$. It should be stressed that the factorization of the Coulomb effects does not take place at $\langle|\vec{x}^*|\rangle$ comparable or larger than the Bohr radius, *e.g.,* in the case when particles are produced in the decays of long lived resonances (η, η') or in the evaporation processes. In such situation, the correlation function tends to 1 and not to the Coulomb factor $A_c(k^*)$ [37].

It can be shown that at $\langle|\vec{x}^*|\rangle > a_c$ and k^* not too close to the classical boundary $(a_c\langle|\vec{x}^*|\rangle)^{-1/2}$, the conditions of the applicability of the classical approach are fulfilled. Thus, *e.g.,* for particles produced in the evaporation processes, the classical trajectory calculations can be used to study the three-body problem of the particle interaction in the Coulomb field of the residual nucleus. Such calculations were recently performed and compared with the two-body quantum ones and the experimental data on pp and pd correlations in the reaction $Ar + Ag$ at 44 MeV/A [41]. The influence of the residual nucleus on the correlation function is found to be of minor importance despite the one-particle spectra are substantially affected. In this context, it would be interesting to select particles in the region of low momenta in which the three body effects should become more important. At present the data from the precise experiment recently performed at GANIL are analyzed to further clarify this problem. Sometimes the effect of the Coulomb field of the residual nucleus is taken into account by the quasi-classical shift of the particle momenta (see, *e.g.*, [42]). Since this shift is typically several tens of MeV/c, it could strongly influence correlations of particles with different mass or electric charge ($pd, pn, \ldots$). Such a procedure seems however incorrect since QS and FSI (coherently including the effect of residual nucleus) affect the dependence of the production amplitude on the final (measured) particle momenta.

The considered approach to FSI cannot be directly applied to the particles produced inside a dense system, *e.g.,* in heavy ion collisions. In this case, the particle interaction with a substantial momentum transfer can be localized, so that we can consider it as the usual scattering, independent of the previous production process. It can be shown that such scattering just reduces to the redistribution of the space-time production points [43]. It means that the interference correlations of the interacting particles are sensitive only to the final production points, distributed near the surface of the dense system. The correlations can be also influenced by the absorption effects [44]. Since the absorption probability is smaller for particles moving towards the nearest boundary of the dense systems, the azimuthal correlations can arise and, as a result, the low-Q correlations can be enhanced. Both these effects have been recently observed [45, 46].

5. Conclusion

The particle interferometry represents an important tool allowing one to study the space-time development of the particle production process (a sort of the space-time "nuclear microscope"). It can be also used as a sensitive measure distinguishing between various models providing the space-time coordinates of the particle production points.

The space-time measurement is however done indirectly, with the help of the correlation function in momentum space. Besides the problems of the reference distribution and the model dependence of the space-time parametrization of the correlation function (*e.g.,* different types of the sources, including resonances), the complications arise due to various effects distorting the correlation function, such as final state interactions, rescattering and absorption. At the same time, these distortions, being dependent on the characteristic distance between the particle production points, also yield the information about the space-time development of the production process. In particular, the nonidentical particle interferometry is possible due to the effect of final state interactions.

It should be stressed that the particle interferometry is of practical interest provided that the characteristic coherence length and time (related to the space-time size of the particle sources) are much smaller than the space and time dimensions of the particle production region, and that the mean energy of the relative motion of the produced multiparticle system is substantial to ensure the validity of the adiabatic approximation in the treatment of the final state interactions (the factorization of the final state interaction between the particles with a small relative velocity).

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