Bose-Einstein Correlations from Random Walk Models

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We argue that strong final state rescattering among the secondary particles created in relativistic heavy ion collisions is essential to understand the measured Bose-Einstein correlations. The recently suggested “random walk models” which contain only initial state scattering are unable to reproduce the measured magnitude and $K_{\perp}$-dependence of $R_{\perp}$ in Pb+Pb collisions and the increase of $R_{t}$ with increasing size of the collision system.

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An important aspect of understanding ultrarelativistic heavy ion collisions is the clear identification of genuine collective nuclear effects which cannot be explained in terms of a simple superposition of nucleon-nucleon collisions. An integral part of our task to look for new physics must therefore be the careful construction of models for nucleus-nucleus collisions (“A-B collisions”) based on a superposition of individual nucleon-nucleon ($N-N$) collisions in order to establish where they fail.

Recently there have been some renewed attempts to construct such models. In the “Random Walk Model” (RWM) of the single-particle transverse mass spectra measured in A-B collision have been calculated by extrapolating those from pA collisions. The LEXUS model of goes even further, by simulating A-B collisions as a simple folding of independent N-N collisions.

Here we demonstrate that, while these models have had some success in describing measured single-particle spectra, they fail to reproduce crucial features of the observed two-particle Bose-Einstein correlations. This is shown to be due to their lack of final state rescattering.

The idea behind the RWM was to provide an alternative interpretation of the measured single-particle $m_{\perp}$ spectra, opposite in spirit to the popular hydrodynamical parametrizations. The latter are based on the assumption of a locally thermalized hadron resonance gas which undergoes longitudinal and transverse hydrodynamic expansion. The transverse expansion is interpreted as a genuine nuclear collective effect with no analogue in N-N collisions; it is identified through the characteristic flattening it causes in the transverse mass spectra, especially at small $M_{\perp} < 2m_{0}$ where the inverse slope parameter is found to increase linearly with the rest mass $m_{0}$ of the produced hadrons.

The RWM, on the other hand, starts from the observation that such a flattening, relative to N-N collisions, happens even in pA collisions where hydrodynamic transverse expansion is not expected to occur. In the RWM an incident nucleon undergoes multiple collisions in the nuclear target, leading to a random walk pattern in the transverse momentum plane. When it collides inelastically it creates a little “fireball”, identical to those formed in elementary inelastic pp collisions, but moving with a transverse rapidity $\rho$ which is Gaussian distributed. In A-B collisions the same mechanism works for both projectile and target nucleons. The width of the $\rho$ distribution is fixed in pA collisions and then extrapolated to A-B collisions using geometric considerations. Resonance decays are neglected, limiting the applicability of the model to sufficiently large transverse momenta.

In LEXUS the same idea is implemented more directly in terms of a folding algorithm for $N-N$ collisions which, in contrast to RWM, also includes longitudinal momentum degradation in multiple collisions. However, the latter is not coupled to the broadening of the $m_{\perp}$-distributions which is parametrized in terms of an effective temperature (inverse slope) which again follows a simple random walk rule.

In both models the transverse broadening parameter is taken as independent of rapidity and not correlated with the position of the collision point (fireball) within the reaction zone. Secondaries do not interact, and (except for energy degradation effects) all primary collisions are alike. For this reason neither model can explain the different chemical composition in pp and A-B collisions (see); they both make statements only about the shape of the momentum distributions.

Existing tests of the models are still very superficial. In the LEXUS output was compared to pion and proton rapidity spectra and midrapidity transverse mass spectra from 200 A GeV S+S collisions at the SPS. In Refs. recent data on single-particle $m_{\perp}$-spectra from 160 A GeV Pb+Pb collisions at the SPS, taken by the NA44 and NA49 collaborations, were compared to the RWM and to hydrodynamical parametrizations. Both types of models give a reasonable description of the data; discrepancies occur mostly at low $p_{\perp}$, where resonance decays distort the pion spectra and hydrodynamical flow effects show up most strongly for the heavier hadrons. In particular, the RWM has difficulties to reproduce the strong rest mass dependence of the transverse slope parameters published by NA44 which have been quoted as evidence for collective transverse flow. It is clear from these papers, however, that without a systematic
study of the $m_\perp$-spectra as a function of the collision system and of rapidity, carefully including resonance decays, single particle spectra alone may not lead to definite conclusions about the validity of the RWM approach. In the present note we therefore examine the discriminating power of two-particle Bose-Einstein correlations.

For ease of language our discussion will be based on the RWM, with comments on LEXUS to follow below. To calculate the correlation function for pairs of identical particles we need to know the emission function $S_{\text{rw}}(x, p)$ describing the phase-space distribution of the particles emitted from the source. Since both the RWM and LEXUS provide explicitly only momentum space information, we must try to reconstruct the corresponding coordinate space information from the description of the models given in Refs. [1,2]. From the emission function the correlator is calculated via [3–5]

$$C_{\text{rw}}(q, K) = 1 + \left| \frac{\int d^4x S_{\text{rw}}(x, K) e^{iq\cdot x}}{\int d^4x S_{\text{rw}}(x, K)} \right|^2,$$  

(1)

where $K = (p_1 + p_2)/2$ is the average pair momentum and $q = p_1 - p_2$ the difference between the two observed on-shell momenta. Since $p_1$ and $p_2$ are on-shell, $q$ and $K$ satisfy the constraint

$q \cdot K = 0 \implies q^0 = q \cdot K^{0} = q \cdot \frac{p_1 + p_2}{E_1 + E_2}$.  

(2)

The single-particle spectrum is calculated as

$$E_{p} \frac{dN}{dp} = P_1(p) = \int d^4x S_{\text{rw}}(x, p, E_{p}).$$  

(3)

For the RWM it was given in the form [1]

$$2\frac{dN}{dy dp_\perp^2 d\phi} = \int d^4x S_{\text{rw}}(x, p) \Phi(\theta(y_L - |Y - Y_0|) \times \frac{V_0 m_\perp}{(2\pi)^3} \cosh(Y - y) \exp\left(-\frac{p \cdot u}{T}\right).$$  

(4)

Here $y, p, \phi$ describe the momentum $p$ of the measured particle, $T$ is the temperature of the fireballs formed in the individual $N$-$N$ collisions, and $Y$ is their longitudinal rapidity which is distributed with a box distribution $\theta(Y_L - |Y - Y_0|)$ between $Y_0 - Y_L$ and $Y_0 + Y_L$. $u$ is the 4-velocity of the fireball, parametrized according to $u(Y, \rho, \Phi) = (\cosh \rho \cos Y, \sinh \rho \cos \Phi, \sin \rho \sin \Phi, \cosh \rho \sin Y) = \gamma(1, u)$ where $\rho$ is the transverse rapidity of the fireball and $\Phi$ its direction. This gives

$$p \cdot u = m_\perp \cosh(Y - y) \cos \rho - p_\perp \cos(\Phi - \phi) \sin \rho.$$  

(5)

The transverse fireball rapidity $\rho$ is distributed with

$$f(\rho) = f_{\text{AB}}(\rho) = \frac{2}{\sqrt{\pi} \delta_{\text{AB}}^2} \exp\left(-\frac{\rho^2}{\delta_{\text{AB}}^2}\right).$$  

(6)

The width of this distribution depends on the collision system $A + B$ via the random walk rule

$$\delta_{AB}^2 = (N_A + N_B - 2) \delta^2,$$  

(7)

where $N_A$ ($N_B$) is the average number of nucleons in nucleus $A$ ($B$) hit by a nucleon from nucleus $B$ ($A$). If the projectile is just one nucleon, $N_A = 1$. The parameter $\delta$ is obtained by fitting $pA$ data ($N_B = 1$).

If one tries to reconstruct $S_{\text{rw}}$ by comparing [3] with [4] one is faced by certain ambiguities, although the appearance of a simple volume factor $V_0$ in front of the integral over the fireball 4-velocity $u$ in [3] excludes most functions $S_{\text{rw}}(x, p)$ with complicated $x$-dependences. The simplest assumption that the longitudinal and transverse fireball rapidities $Y$ and $\rho$ of the fireballs are not correlated with the fireball positions is, however, untenable. In this case the emission function factorizes, $S_{\text{rw}}(x, p) = F(x) \cdot I(p)$ (with $I(p)$ given by the r.h.s. of Eq. [1] divided by $V_0$), and [3] yields a correlator which does not depend on $K$. Since all heavy-ion data show a clear dependence of the correlation function (in particular of its longitudinal width parameter $R_L$) on the pair momentum $K$ [1,2,3], such an emission function is excluded.

More realistically one should at least implement the simple expectation that (in RWM language) fast fireballs live longer and fly farther before decaying than slow ones. This leads to a correlation between the decaying fireball’s position $x$ and velocity $u$, inducing also a correlation between $x$ and $p$ via the Boltzmann factor and thereby a $K$-dependence of the correlator. We thus write

$$S_{\text{rw}}(x, p) = \int d^4x d\rho d\Phi F(x, Y, \rho) \times \frac{m_\perp}{(2\pi)^3} \cosh(Y - y) \exp\left(-\frac{p \cdot u}{T}\right),$$  

(8)

with $\int d^4x F(x, Y, \rho) = V_0 \theta(Y_L - |Y - Y_0|)$. Azimuthal symmetry excludes a $\Phi$-dependence of $F$. We will use longitudinally boost-invariant space-time coordinates $x^\mu = (t, \eta, r, \phi)$, with $\tau^2 = t^2 - z^2$, $\eta = \frac{1}{2} \ln[(t + z)/(t - z)]$, and integration measure $d^4x = \tau dr d\eta d\phi$. In the RWM [1] the transverse fireball rapidity is not correlated with its space-time position since the width $\delta_{AB}$ of $f(\rho)$ is determined from the impact parameter averaged number of hit nucleons. Thus there are no transverse $x$-$p$ correlations in the source. This is easily seen to lead to a correlator whose transverse width, given by the transverse HBT radius $R_T$, does not depend on $K$. This contradicts the experimental observation of a clear and rather strong decrease of $R_T$ as a function of $K_{\perp}$ in 160 A GeV Pb+Pb collisions [11] which cannot be explained by resonance decays [13,14].

One can try to include the missing transverse $x$-$p$ correlations by making the broadening parameter $\delta_{AB}$ r-dependent. Instead of using the impact parameter averaged number of hit nucleons from the original formulation of the RWM [1,2,3],

$$N_A = (2/3) \pi R_{A}^2 2 R_A n_0,$$  

(9)
(where \( r_0 = 0.8 \text{ fm} \) is the nucleon radius, \( R_A = 1.12 A^{1/3} \text{ fm} \) is the nuclear radius, and \( n_0 = 0.17 \text{ fm}^{-3} \) is the standard nuclear density), we can take the actual value at distance \( r \) from the collision axis:

\[
N_A(r) = (\pi r_0^2) 2 \sqrt{R_A^2 - r^2} n_0 .
\]

(10)

Via Eq. (9) this yields an \( r \)-dependent width \( \delta_{AB} \) and, via (3), an \( r \)-dependent \( \rho \)-distribution \( f_{AB}(\rho, r) \). We will now study whether the resulting transverse \( x \)-\( p \) correlations can cause the measured \( K_{1 \perp} \)-dependence of \( R_{\perp} \).

Since the RWM does not take into account longitudinal energy degradation by multiple initial state scattering, we can factorize the \( \rho \) and \( Y \) dependence of the function \( F \) in (8): \( F(x; Y, \rho) = F'(x; Y) f(\rho, r) \). There is no correlation between the transverse fireball rapidity and its longitudinal position. Furthermore, without longitudinal momentum loss the maximum and minimum longitudinal fireball rapidities \( Y_0 \pm Y_L \) are not related to the actual number of hit nucleons and thus independent of \( r \). Hence \( F'(x; Y) \) can be further factored into \( H(\eta, \tau; Y) \cdot G(r, \varphi) \). The transverse HBT radius \( R_{\perp} \) can be calculated from [17]

\[
R_{\perp}^2 = \langle r^2 \sin^2 \varphi \rangle, \quad \langle f(x) \rangle = \frac{\int d^4x f(x) S(x, K)}{\int d^4x S(x, K)} .
\]

(11)

Since the radial integral couples \( G \) and \( f \), we must know the functional form of \( G(r, \varphi) \). We take \( G \) proportional to the number of collisions at distance \( r \) from the collision axis; for nuclear collisions at zero impact parameter this is given by

\[
G(r, \varphi) = G(r) \propto \sqrt{R_A^2 - r^2} \sqrt{R_B^2 - r^2} .
\]

(12)

In the following we consider only symmetric collisions, \( R_A = R_B \). The generalization to the case \( R_A \neq R_B \) is obvious.

The evaluation of (11) is complicated by the fact that Eq. (3) couples \( r \) to \( Y \) such that the function \( H(\eta, \tau; Y) \) must also be known. To obtain an as much as possible model independent estimate for the behaviour of \( R_{\perp} \) we use the following simplification: We go to the LCMS (Longitudinally Co-Moving System, \( y = 0 \)) and fix \( Y = 0 \), thus taking into account pion production only from fireballs which rest in the LCMS. Since contributions from other fireballs are suppressed by a factor \( \exp[-M_{\perp}(\cosh Y - 1) \cosh \rho/T] \), we expect our simplification to provide a good estimate for the correct \( R_{\perp} \). The advantage of this approximation is that now the integration in Eq. (11) factorizes in \( (\tau, \eta) \) and \( (r, \varphi) \). One finds

\[
R_{\perp}^2 \approx \left\{ \int d\rho \int r dr r^2 C(r, \rho) \right\} / \left\{ \int d\rho \int r dr 2 C(r, \rho) \right\} \approx \left\{ \int r dr r^2 C(r, \rho) \right\} / \left\{ \int r dr 2 C(r, \rho) \right\} .
\]

(13)

where

\[
C(r, \rho) = (R_A^2 - r^2) \left( \pi \left( 4\pi r_0^2 n_0 \sqrt{R_A^2 - r^2} - 2 \right) \delta^2 \right)^{-1/2}
\]

\[
\times \exp \left[ \frac{-\rho^2}{(4\pi r_0^2 n_0 \sqrt{R_A^2 - r^2} - 2)\delta^2} \right]
\]

\[
\times \exp \left( -\frac{M_{\perp}}{T} \cosh \rho \right) I_0 \left( \frac{K_{\perp}}{T} \sinh \rho \right) .
\]

(14)

Here the integrations over \( \Phi \) and \( \varphi \) have already been performed. Note that the \( r \)-integration in Eq. (13) does not extend to \( R_A \) but only to \( R_A - \delta \) where \( \delta \) is chosen such that at transverse distance \( r = R_A - \delta \) there is still one nucleon-nucleon collision. This is ensured by requiring \( \delta_{AA}(R_A - \delta) = 0 \); if \( r > R_A - \delta \) then \( \delta_{AA}(r) \) is negative due to Eqs. (11) and (10).

![FIG. 1. The \( M_{\perp} \)-dependence of \( R_{\perp} \) according to (13) for Pb+Pb (A=207) and S+S (A=32) collisions.](#)

For our computation we re-calculated \( \delta \) according to Eq. (9) (see (10)). Using \( \delta_{\text{MW}} = 0.3 \) we got \( \delta = 0.217 \). Results for \( R_{\perp} \) according to the approximation (13) for S+S (A=32) and Pb+Pb (A=207) reactions are shown in Fig. 1. It is immediately clear that the calculated values for \( R_{\perp} \) underestimate the experimental data (which, at small \( M_{\perp} \), give \( R_{\perp} \approx 6 \text{ fm} \) for Pb+Pb [11] and \( \approx 4 \text{ fm} \) for S+S [12,13]) considerably. It has been argued in [18] that the present data on \( R_{\perp} \) from Pb+Pb collisions require a strong transverse expansion since the measured values of \( R_{\perp} \) signal a source much bigger than the original Pb nucleus. The \( R_{\perp} \)-values in Fig. 1, however, are consistent with a transversally non-expanding source. Furthermore, the \( M_{\perp} \)-dependence of \( R_{\perp} \) induced by the \( r \)-dependence of \( \delta_{AB} \) is much too weak and cannot reproduce the strong \( M_{\perp} \)-dependence seen in the Pb+Pb data [11].

The wrong magnitude and \( K \)-independence of \( R_{\perp} \) are, however, not the only problems of the RWM/LEXUS models. Following [19] we now argue that they are also unable to correctly describe the dependence of the longitudinal radius parameter \( R_l \) on the size of the collision system.
In RWM and LEXUS the full source of secondary particles consists of smaller sources generated in individual N-N collisions. Let \( s_c(x, p; x_0) \) be the emission function describing such a small source where \( x_0 \) stands for the point where the N-N collision occurs. Translation invariance gives \( s_c(x, p; x_0) = s_c(x - x_0, p) \), and the total emission function can be written as

\[
S(x, p) = \int d^3x_0 \sigma(x_0) s_c(x - x_0, p),
\]

(15)

where \( \sigma(x_0) \) is the distribution of collision points resulting from the collision geometry. Its normalization \( \int d^4x_0 \sigma(x_0) \) is fixed by the number of N-N collisions.

We are now interested in the \( q_1 \)-dependence of the correlator. Inserting (13) into (4), setting \( q_1 = 0 \), and boosting into the LCMS \((\beta_l = 0 = K_l)\) we find

\[
C(q, K) \bigg|_{q_1 = K_1 = 0} = \frac{\int d^4x_0 d^4x_0 e^{-iqz} \sigma(x_0) s_c(x - x_0, K)^2}{\int d^4x_0 d^4x_0 \sigma(x_0) s_c(x - x_0, K)^2} \left( \int d^4x e^{-iqz} \sigma(x)^2 \right)^2
\]

(16)

\[
= \frac{\int d^4x d^4x e^{-iqz} s_c(x, K)^2}{\int d^4x s_c(x, K)} \left( \int d^4x e^{-iqz} \sigma(x)^2 \right)^2
\]

(17)

where \( z = \tau \sinh \eta \). This expression is typical for LEXUS type models where \( \sigma(x) \) is determined by collision geometry only. However, an additional \( K \)-dependence of \( \sigma \) would not affect the following argument. The correlator would still factorize, and \( \sigma(x) \) will still be non-zero only in the region where N-N collisions occur.

The correlator (17) is in fact a product of two correlators associated with the sources \( s_c \) and \( \sigma \), respectively. The longitudinal radius parameter of the whole source extracted from a Gaussian fit to (17) is thus given by

\[
R_l^2 = R_{l,s_c}^2 + R_{l,\sigma}^2,
\]

(18)

where \( R_{l,s_c} \) \((R_{l,\sigma})\) is the longitudinal radius parameter associated with the source function \( s_c \) \((\sigma)\). \( R_{l,s_c} \) can be extracted from correlation measurements in N-N collisions. Based on data published in [23] this value was estimated in [24] to be \( R_{l,s_c} \approx 1.7 \) fm.

When considering different collision systems clearly only \( R_{l,\sigma} \) changes. It can be calculated using the model-independent expression [23]

\[
R_{l,\sigma}^2 = \langle z^2 \rangle_\sigma - \langle z \rangle_\sigma^2
\]

(19)

where \( \langle f(x) \rangle_\sigma = \int d^4x f(x) \sigma(x) / \int d^4x \sigma(x) \).

The compilation in Fig. 1 of Ref. [19] shows that in A-B collisions \( R_{l,\sigma} \) rises approximately linearly with \((A - 1)^{1/3} + (B - 1)^{1/3} \), with a slope which is close to 0.5. The geometric distributions \( \sigma(x) \) for the collision points in RWM or LEXUS cannot (via Eq. (13)) reproduce such a strong rise [19]. This can be seen immediately by calculating the r.h.s. of (17) in an approximation where the two nuclei are replaced by cylinders with the same volume and \( \langle z^2 \rangle \). One finds

\[
R_{l,\sigma}^2 = \frac{1}{20} \frac{1}{\gamma^2} (R_A^2 + R_B^2),
\]

(20)

where \( \gamma \approx 10 \) is the Lorentz contraction factor in the N-N CMS at CERN/SPS energies. The factor \( 1/(20 \gamma^2) \) makes \( R_{l,\sigma} \) rise by about a factor 6 more slowly than required by the data; for Pb+Pb collisions, after adding the contribution \( R_{l,s_c}^2 \approx 3 \) fm\(^2\), Eqs. (16,20) underpredict the data [11] for \( R_{l,\sigma}^2 \) by a factor 10. Hydrodynamic parametrizations [23][24] of the source function, on the other hand, can easily accommodate this rise by adjusting the lifetime \( \tau_0 \) of the collision fireball. For Pb+Pb collisions at CERN one finds \( \tau_0 \approx 8 \) fm/c [23]. This parameter measures the length of time during which the secondary particles created in the nuclear collision rescatter among each other. In scenarios with approximately boost-invariant longitudinal expansion \( R_l \) is directly proportional to this time. Models without rescattering (like RWM and LEXUS) do not lead to such a collective longitudinal expansion and thus cannot reproduce \( R_{l,\sigma} \).

In summary we conclude that the RWM and LEXUS models, which try to describe the hadronic momentum spectra in A-B collisions in terms of a linear extrapolation of pA respectively pp spectra, cannot reproduce the magnitude of the transverse radius parameter \( R_{l,\sigma} \) and its dependence on the transverse momentum nor the dependence of the longitudinal radius parameter \( R_l \) on the size of the collision system. The first failure is due to the lack of transverse expansion and transverse space-momentum correlations in these models. The second failure is caused by the lack of rescattering which would result in longer lifetimes and larger values of \( R_{l,\sigma} \) through collective longitudinal expansion.

Simple improvements of these models which preserve their original spirit as a superposition of N-N collisions cannot remedy these failures. Extensive final state rescattering among the produced hadrons, as implemented in other Monte Carlo transport models and in hydrodynamical source parametrizations, is essential to reproduce the correlation data.

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