Emission angle dependent HBT at RHIC and beyond

Peter F. Kolb\textsuperscript{a} and Ulrich Heinz\textsuperscript{b}

\textsuperscript{a}Department of Physics and Astronomy, SUNY at Stony Brook, NY 11794-3800, USA
\textsuperscript{b}Physics Department, The Ohio State University, Columbus, OH 43210, USA

We study the geometrical features of non-central heavy ion collisions throughout their dynamical evolution from equilibration to thermal freeze-out within a hydrodynamic picture. We discuss resulting observables, in particular the emission angle dependence of the HBT radii and the relation of these oscillations to the geometry at the final stage.

1. Introduction

The systematic investigation of observables generated by the broken azimuthal symmetry of non-central nuclear collisions \cite{1} has lead to major insight in the dynamics of the expansion. In particular the large signal of anisotropic particle flow \cite{2} requires a surprisingly efficient microscopic rescattering mechanisms \cite{3} and even agrees with the limit of infinitely short rescattering lengths, namely hydrodynamic calculations \cite{4,5}, thus providing access to the nuclear equation of state under extreme conditions. Furthermore, the large signal is critically dependent on rapid thermalization and early pressure in the system, at timescales smaller than 1 fm/c \cite{7}. Experimentally and theoretically even more challenging is the extraction and interpretation of the azimuthal dependence of coordinate space observables such as two-particle correlations \cite{8,9}. Its understanding however offers an additional handle to study the space-time geometry of the source, helps to understand and clarify persisting problems \cite{12}, and ultimately is one more piece in our effort to understand the \textit{complete} dynamics of heavy ion collision.

2. Fireball initialization and evolution

To elucidate the connection of these observables and the geometry that they reflect, we study two scenarios in the following. In the first one (referred to as ‘RHIC1’) we choose the initial and freeze-out conditions such that the particle yield as well as the pion and antiproton spectral slopes agree well with experimental data from \textit{central} Au+Au collisions at 130 AGeV (for details, see \cite{4}). The maximum temperature in the center of the collision region is found to be $T_0 = 340$ MeV at an assumed equilibration time of $\tau_0 = 0.6$ fm/c. The spectral slopes require a decoupling temperature of $T_{\text{dec}} = 130$ MeV, where fluid elements are assumed to liberate their particle content boosted by the local flow velocity. \textit{Geometry}, without involving further parameters, then determines the initialization of \textit{non-central} collisions. Experimental results on elliptic flow confirmed the
large values expected from hydrodynamic predictions \cite{7} and justified this idealized approach as long as the produced fireball is sufficiently large to allow for thermalization and ‘macroscopic behavior’, a requirement which expectedly breaks down above a certain impact parameter. For Au+Au collisions at $\sqrt{s_{NN}} = 130$ GeV the data indicate that this limit is reached for $b \sim 7-8$ fm \cite{4}.

With the second scenario we want to stress the geometrical features of the source and the imprints on experimental observables. To enhance the relevant signals, we employ more extreme conditions of $T_0 = 2$ GeV at $\tau_0 = 0.1$ fm/c and $T_{\text{dec}} = 100$ MeV. As the system stays above the QCD transition temperature for most of the time of the evolution we employ an ideal gas equation of state for simplicity (for reasons that will become clear later we label this system ‘IPES’). At the moment it is not clear whether these conditions are accessible at future colliders, but the calculation allows us to trace back the origin of qualitative changes in the observables which might already occur at LHC energies. In the following we study the evolution of both systems for an impact parameter of 7 fm on the basis of an ideal hydrodynamic evolution with longitudinal boost-invariance \cite{7}.

Upon impact the nuclei produce secondaries and large amounts of entropy in their geometrical overlap region. This is generally modeled by contributions from wounded nucleons and/or binary collisions to the initial energy- or entropy density fields in the overlap region \cite{4}. We characterize the spatial eccentricity of the fields by averaging over the transverse plane (at $z = 0$) with the energy density as weight. Defining $\epsilon_x = \langle y^2 - x^2 \rangle / \langle y^2 + x^2 \rangle$, we can follow its evolution through time (see Fig. 1). If the impact parameter points in $x$-direction, $\epsilon_x$ is positive during the initial stages when the overlap-ellipsoid is elongated out of the reaction plane. As time evolves, the stronger pressure gradients in the short direction force a stronger expansion in plane than out of plane, thereby reducing the initial anisotropy and eventually leading to an overshoot such that the source appears in plane elongated at the very end of its existence ($\epsilon_x < 0$). This is in particular true for the ‘IPES’ case, which exhibits much stronger pressure gradients and therefore stronger transverse expansion (‘IPES’ stands for ‘In Plane Elongated Source’). The generation of the accompanying anisotropic flow is characterized through the momentum anisotropy $\epsilon_p = \langle T_{xx} - T_{yy} \rangle / \langle T_{xx} + T_{yy} \rangle$ where the $T^{ij}$ are the diagonal elements of the energy mo-

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig1.png}
\caption{Time evolution of the spatial eccentricity $\epsilon_x$ and momentum anisotropy $\epsilon_p$ for RHIC1 (solid) and IPES (dashed).}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig2.png}
\caption{Freeze-out hypersurface of the two systems. Shown are cuts at $y = 0$ (solid) and $x = 0$ (dashed).}
\end{figure}
mentum tensor of the fluid. Ultimately, this momentum anisotropy is reflected as elliptic flow. Fig. 1 shows clearly how the geometrical eccentricity causes a rapid buildup of momentum anisotropy (the early saturation of $\epsilon_p$ at about 3 fm/c in the RHIC1 case is caused by the mixed phase of the employed equation of state [7]).

As particle emission occurs throughout the evolution from the surface of the fireball, it is not clear what average geometry will be imprinted on the particle distributions. Particles with high transverse momenta preferentially emerge from regions with the largest collective transverse velocities, that is from the freeze-out hypersurface at its largest radial extension, whereas the low $p_\perp$ particles are emitted when the interior of the fireball freezes out in a timelike manner (see the freeze-out hypersurfaces in Fig. 2). Therefore particles of different $p_\perp$ (or particle pairs with different mean transverse momentum $K_\perp$ as they are studied in HBT) probe the system at different times and different regions of the fireball.

3. Azimuthally sensitive HBT-analysis

From the hydrodynamically determined freeze-out hypersurface $\Sigma$ we evaluate the density of emitted pions of momentum $K$ through $S(x, K) \propto \int_\Sigma K \cdot d^3\sigma(x') f_{BE}(E(x'), T_{dec}) \times \delta^4(x - x')$ with $E(x') = K \cdot u(x')$ where $u$ is the local flow velocity. With this source function $S(x, K)$ as a weight we can calculate the average emission points and variances and collect them in the spatial correlation tensor $S_{\mu\nu} = \langle x_\mu x_\nu \rangle - \langle x_\mu \rangle \langle x_\nu \rangle$. Rotating this correlation tensor into the direction of the particle pair momentum (angle $\Phi$), one obtains the ‘width’ and ‘depth’ of the emission region with respect to the direction of observation and can make the connection [9] to the experimentally deduced HBT radii [8].

The ‘sideward’ radius receives contributions only from the $(x, y)$-components of the correlation tensor and thus can be related directly to the geometry of the fireball. More specifically it reflects the ‘width’ of the emission region of particles of similar momentum orthogonal to their average momentum (the direction of observation). For large transverse momenta, these ‘homogeneity regions’ are narrow slivers close to the rim of the freeze-out hypersurface whose shape is quite different from that of the total source [10]. The ‘out’ and ‘out-side’ terms mix geometrical and temporal emission information, with the temporal part gaining weight with increasing transverse momentum. In the calculations we can investigate both parts separately, but experimentally they appear always in combination.

Fig. 3 displays the oscillations of the squared radii for the RHIC1 calculations. Each radius shows its own oscillation pattern which remains qualitatively unchanged for all transverse momenta. The sideward oscillation reflects the out-of-plane extended source: observing such a source from the $x$-direction it appears to be larger sideways than when observed from the $y$-direction. This remains true even at large $K_\perp$ where the emission is concentrated along a thin sliver near the rim of the fireball [10]. The outward oscillations confirm this behavior from an ‘orthogonal’ viewpoint. The sign of the out-side oscillation can be traced back to a ‘tilt’ of the emission region relative to the emission direction [10] which is generated by the out-of-plane deformation of the source. Even though hydrodynamic calculations fail in the quantitative description of the freeze-out geometry [12], the overall source shape appears to be in agreement with preliminary data [8].

Moving to the IPES calculation a drastic change occurs (Fig. 4): The oscillations flip sign at $K_\perp \sim 0.15$ GeV, showing the same oscillation pattern as at RHIC1 for larger
$K_\perp$ but opposite oscillations below. The geometric contributions to the out and out-side radius (marked by circles) always oscillate opposite to the RHIC1 case, reflecting the in-plane elongation of this source. At larger $K_\perp$ this naively expected behavior is masked by growing temporal contributions, leading eventually to the same qualitative oscillation pattern as seen for RHIC1. Because the homogeneity regions do not trace out the interior of the source but concentrate near the rim of the system, the oscillations of the sideward radius also change sign at larger $K_\perp$. The true geometry of the source is therefore only reflected in the correlations of the particles with the smallest transverse momenta.

Acknowledgements: This work was supported in part by the U.S. Department of Energy under Contracts No. DE-FG02-88ER40388 and DE-FG02-01ER41190. PFK is supported by a Feodor Lynen fellowship of the Alexander von Humboldt Foundation.

REFERENCES

1. See the contributions of the RHIC collaborations, these proceedings.
2. J.-Y. Ollitrault, Phys. Rev. D 46 (1992) 229.
3. D. Molnar and M. Gyulassy, Nucl. Phys. A 697 (2002) 495, [Erratum ibid. A 703 (2002) 893]; Z.W. Lin and C.M. Ko, Phys. Rev. C 65 (2002) 034904.
4. P.F. Kolb et al, Phys. Lett. B 500 (2001) 232; P. Huovinen et al., Phys. Lett. B 503 (2001) 58; P.F. Kolb et al., Nucl. Phys. A 696 (2001) 197.
5. D. Teaney, J. Lauret, and E.V. Shuryak, Phys. Rev. Lett. 86 (2001) 4783; nucl-th/0110037
6. T. Hirano, Phys. Rev. C 65 (2002) 011901.
7. P.F. Kolb, J. Sollfrank, and U. Heinz, Phys. Lett. B 459 (1999) 667; Phys. Rev. C 62 (2000) 054909.
8. F. Retière et al. (STAR Collab.), nucl-ex/0111013; M. López-Noriega, these proceedings.
9. U.A. Wiedemann, Phys. Rev. C 57 (1998) 266.
10. U. Heinz and P.F. Kolb, Phys. Lett. B 542 (2002) 216 [hep-ph/0206278].
11. U. Heinz, A. Hummel, M.A. Lisa, and U.A. Wiedemann, nucl-th/0207003
12. U. Heinz and P.F. Kolb, Nucl. Phys. A 702 (2002) 269; hep-ph/0204061.