INTERFEROMETRY OF HIGH ENERGY NUCLEAR COLLISIONS

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Interferometry or HBT effect or GGLP effect refers to the same phenomenon, corresponding to the possibility of determining space-time dimensions of particle emitting sources. This is accomplished by the width of the second order correlation function, the term "second" referring to intensity interference as opposed to the commonly known amplitude interference in Optics. It was discovered in the mid fifties by Hanbury-Brown and Twiss[1], after whom it was named. Later, in 1960, analogous effect was observed by Goldhaber et al.[2] when studying correlations between equal pions produced in $\bar{p}p$ collisions. After that, it was also known as GGLP effect. More recently, along the last decade, the phenomenon was faced as a potential powerful tool for probing the Quark Gluon Plasma (QGP), expected to be formed in high energy nuclear collisions or in high multiplicity fluctuations in $\bar{p}p$ collisions. Interferometry then would be useful for measuring the expected large space-time dimensions of the QGP.

For pedagogical purposes, it is useful to start with the case of two point sources for illustrating the concept of interferometry (for interesting reviews, see Ref. [3,4]). Consider then that two point sources I and II located in x_1^μ and x_2^μ emit identical quanta which are simultaneously observed, with momenta k_1^μ and k_2^μ, at detectors A and B, located at x_A^μ and x_B^μ, respectively, as illustrate in Figure 1.

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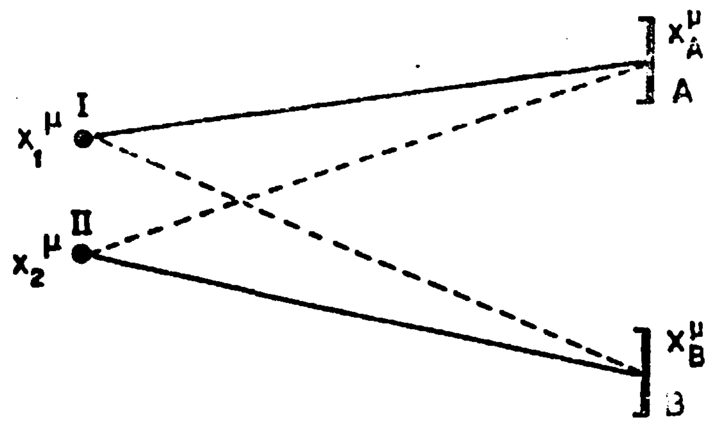

FIGURE 1:

Since the two particle are identical, there are two possibilities for writing the amplitude for simultaneous detection: A detects the particle emitted by source I and B the one emitted by II or vice-versa, i.e.,

$$\begin{aligned} A(k_1, k_2) &= \frac{1}{\sqrt{2}}[\epsilon^{-ik_1.(x_A - x_1)}\epsilon^{i\phi_1}\epsilon^{-ik_2.(x_B - x_2)}\epsilon^{i\phi_2} \\ &\pm \epsilon^{-ik_1.(x_A - x_2)}\epsilon^{i\phi_2'}\epsilon^{-ik_2.(x_B - x_1)}\epsilon^{i\phi_1'}] \ . \end{aligned} \quad (1)$$

where ϕ is the associated emission phase, here assumed independent of k. The (+) sign refers to bosons and the (-) one, to fermions. We are implicitly considering that the equal quanta are pions (unless otherwise stated). The two sources may be chaotic, like a thermal source, or coherent, like a laser. Let us consider here that they are completely chaotic. If so, the phases change in each emission and we should consider an average over phases when calculating the corresponding probability. If we take this average over a time long compared to the phase fluctuations, the probability for observing the two pions simultaneously is given by

$$\begin{aligned} P(k_1, k_2) &= \tfrac{1}{2}[2 \pm (\epsilon^{i(k_1 - k_2).(x_1 - x_2)}\langle \epsilon^{i(\phi_1 + \phi_2 - \phi_1' - \phi_2')}\rangle + c.c.)] \\ &= 1 \pm cos[(k_1 - k_2).(x_1 - x_2)] \ . \end{aligned} \quad (2)$$

The final result follows from the chaoticity assumption and from

$$\langle \epsilon^{\pm i(\phi_1 + \phi_2 - \phi_1' - \phi_2')}\rangle = \delta_{\phi_1\phi_1'}\delta_{\phi_2\phi_2'} + \delta_{\phi_1\phi_2'}\delta_{\phi_2\phi_1'} \ .$$

The two-particle correlation function is then given by

$$C(k_1,k_2)=\frac{P(k_1,k_2)}{P(k_1)P(k_2)}=1\pm cos[(k_1-k_2).(x_1-x_2)] \ , \qquad (3)$$

where $P(k_i)$ are the single inclusive distribution which, for the case of two chaotic point sources equals unity after taking the average over phases during a very long time.

We should notice that, in the case of identical bosons, there is an enhancement of the correlation function when their relative momenta are small compared to the inverse of the typical spatial dimensions of the reaction volume. For large values of their relative momenta, the correlation function should tend to one, what is clearly not the case for (3), but this is just an artifact of having only two point sources. An important fact that should be emphasized is that the HBT/GGLP effect follows not only from the Bose-Einstein symmetrization but also from the fact that the source is incoherent. This means that, for completely chaotic sources, $C(k_1, k_2=0)=2$ and that, for completely coherent sources, $C(k_1, k_2)=1$ for all values of the momentum difference.

Above, we have considered the simple case of only two point sources. More generally, for extended sources with normalized space-time distributions $\rho(x)$, the two-particle inclusive distribution for completely chaotic sources, is written as

$$\begin{aligned} P(k_1,k_2) &= P(k_1)P(k_2)\int d^4x_1d^4x_2 \mid A(k_1,k_2)\mid^2 \rho(x_1)\rho(x_2) \\ &= P(k_1)P(k_2)\{1\pm \mid \rho(q)\mid^2\} \ , \qquad (4) \end{aligned}$$

the second term corresponding to the Fourier transform of $\rho(x)$.

The correlation function, according to (3), is written as

$$C(k_1,k_2)=\{1\pm\lambda\mid\rho(\mathbf{q})\mid^2\} \ , \qquad (5)$$

where λ is the incoherence or chaoticity parameter, introduced by Deutschmann et al[5] for reducing systematic errors when fitting the theoretical curves with experimental results. This was due to the fact that the experimental results hardly reached the maximum value of 2, while the theoretical curves assumed this value in the zero of the variable in terms of which the correlation function was plotted. According to their suggestion then, completely chaotic sources would result in $\lambda = 1$ while completely coherent ones would correspond to $\lambda = 0$. However, we should be very careful when establishing this correspondence in a straightforward way: usually, the correlation function is studied in terms of one of the so-called Kopylov variables, which are defined in Figure 2, e.g., $C(q_T)$ as a function of q_T, while q_L was ideally considered to be zero. In real life, however, we deal with the problem of finite (and sometimes low) statistics for the events. What happens is that q_L is not zero but is considered in a finite interval. This fact, by itself, is enough for making $C(q_T) < 2$ even for completely chaotic sources! This important point was already emphasized in Ref.[22,7].

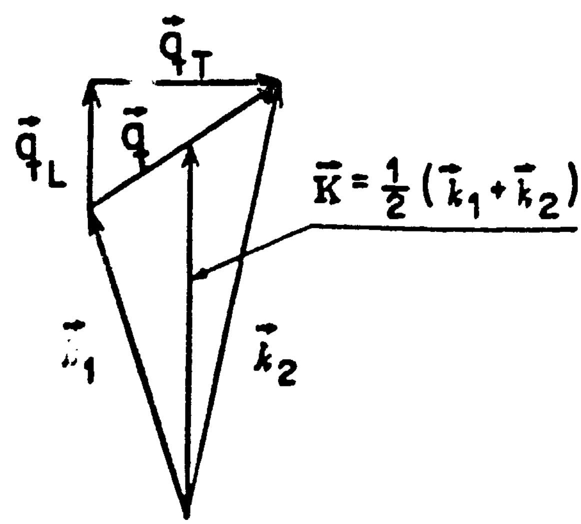

FIGURE 2:

The Kopylov variables q_L and q_T are defined as the projections of the momentum difference $q = k_1 - k_2$, respectively, parallel and perpendicular to the average momentum of the pair, $K = \frac{1}{2}(k_1 + k_2)$.

When dealing with expanding sources, as is the case of high energy collisions, q_L is usually considered as the component of the momentum difference parallel to the collision axis and q_T as the component perpendicular to that direction.

In (5), we see that the correlation function is related to the Fourier transform of the source space-time distribution. Being so, given a certain distribution according a specific model. eq.(5) would return the parameters like radius R and lifetime of the source or the chaoticity parameter, when fitted to the experimental data. However, we should emphasize that this straightforward correspondence between correlation function and geometrical information about the source is only valid when the phase space distribution is completely decoupled, i.e., when $D(x,p) = \rho(x)f(p)$. This is true for the original HBT effect but is usually not valid for high energy collisions, where the source expansion and other dynamical effects may blur the behavior of the correlation function as we will see later. Before discussing this point in more detail, let us open a space for mentioning some of the most significant contributions to the field besides those already mentioned and the more commonly used parametrizations for fitting data.

One of the most popular parametrizations along these last thirty years is the Gaussian or GGLP type, where the space-time distribution of the pion emitting source is considered to be a Gaussian in shape, which results, from (5), into

$$C(k_1,k_2) = 1 + e^{-R^2(\mathbf{q}^2 - q_0^2)} \quad , \tag{6}$$

where R is a parameter that represents an average over space and time dimensions. However, as noted by Bowler[8], this parametrization is good only if $\rho(q)$ contains no directional dependence, i.e., if it depends only on the modulus of the 4-vector q (besides, of course, the phase space decoupling already mentioned above).

Another well known parametrization, proposed by Kopylov. Podgoretskiĭ and Grishin[9] and used in several cases from the beginning to the mid 70's, considers excited systems as sources of pions, emitting them from the surface of a sphere of radius R, with exponential lifetime. As a consequence, the correlation function is written as

$$C(k_1,k_2) = 1 + [\frac{2J_1(q_T R)}{q_T R}]^2 [1 + (q_0 \tau) 2]^{-1} \quad , \tag{7}$$

where $q_0 = E_1 - E_2 \approx q_L$ for small values of the squared (q^2) 4-momentum. In '74 Cocconi[10] re interpreted the quantity $c\tau = \delta(< R)$ as the thickness of the pion emission layer.

We should notice that both parametrizations mentioned above should be restricted to applications involving static sources, since no effect of expansion is being taken into account in their prescription. In what follows, we will just briefly comment about the other models, without mentioning their specific form for the correlation function.

A more general formulation for chaotic sources was suggested in '73 by Shuryak[11], in which any type of source could be introduced. At the end of that decade, on '77, Grassberger[12] pointed out, for the first time, the important rôle that resonances would play in interferometry. This is because, a resonance of 4-momentum k, mass M and width Γ would travel a distance $\sim \frac{k}{M\Gamma}$ before decaying, causing interfering effects whenever $k^{\mu}q_{\mu} \leq M\Gamma$. We will see later on that this effect really seems to be quite significant.

We mentioned above the importance of the source degree of chaoticity for the constructive interference of identical bosons. In '77 and '79, Fowler and Weiner[13] and in '79, Gyulassy, Kauffmann and Wilson citegyu proposed different formalisms for studying the cases of coherent, partial coherent and chaotic sources. They also suggested different approaches for treating final state interactions among the emitted particles, an effect which also may distort the behavior of the correlation function. In this context, more recently, Suzuki[15] in '87 and Bowler[16] in '88, analyzed the distortions introduced into the correlation function coming from s-wave $\pi\pi$ scattering phase shifts in the channel of isospin $I = 2$. Another point we mentioned in connection with eq.(5) was its validity restricted to having uncorrelated phase space (for a detailed discutions of its implications, see Ref.[17]), i.e., space-time distribution decoupled the energy-momentum one. If that was the general case, the formulation proposed by Yano and Koonin[18] in '78 would apply. However, this would not be expected to be the case for pions produced in e^+e^- annihilation. In '85 Bowler[8], as well as Andersson and Hofmann[19] in '86, proposed that the sources in that case would be given by string fragmentation models.

Considering pion interferometry as a powerful probe of the expected quark gluon plasma formation, some models were proposed that had this hypothesis as starting point, i.e., the pions observed in coincidence experiments would have been originated in the hadronization of the plasma. The models by Pratt[20] ('84 - '86), Hama and myself[22] ('86 - '88) and of the group Makhlin, Sinyukov and Averchenkov[23] ('87,'88) considered a hydrodynamic evolution for the plasma. On the other hand, the model by Bertsch, Gong and Tohyama[24], the hadronization happened by evaporation of plasma droplets (globs). The first three groups observed a characteristic symptom of the breakdown of the straightforward geometrical interpretation associated to eq.(5): the correlation function, besides depending on $q = k_1 - k_2$, turned to depend also on $K = \frac{1}{2}(k_1 + k_2)$. The two first groups above also observed that fluctuations in the breakup time would distort the correlation function and that, as a consequence, the $C(\mathbf{q}_T \parallel \mathbf{K}_T)$ would be narrower than $C(\mathbf{q}_T \perp \mathbf{K}_T)$. Later, Bertsch[24] named these variables respectively as $\mathbf{q}_{OUT}$ and $\mathbf{q}_{SID}$ and suggested as a unequivocal way of testing the QGP formation. However, this was shown not to be the case[7], due to the limited statistics still available. But, a necessary test for completely trusting the QGP based models, is to discard more conventional explanations. With this in mind, a resonance gas model based on the ATTILA version of LUND[25] model and treated in terms of the Covariant Current Ensemble[14] with phase-space correlations implicity bilt in, was proposed[7], where the experimental cuts were taken into account in a Monte Carlo code. Actually, as we see in Figure 3, the experimental uncertainties do not allow for favoring any of these classes of models just discussed: all them seem to be compatible with data. Within the same context, Csörgo, Zimányi, Bondorf Heiselberg and Pratt[26] developed afterwards an interferometry code with output coming from SPACER[27] version of LUND model, where resonances were included, as well as experimental cuts. However, they applied a formulation of Yano-Koonin type, which, as already mentioned before, is rigorous only when phase-space correlaitons are weak.

On resuming our discussion about the restrictions imposed on the applicability of eq.(5), we emphasize again that it is only valid if the freeze-out space-time and momentum coordinates of the pions are uncorrelated. In high energy hadronic processes there are many potential sources of such correlations which can significantly modify[20,21,22,23] the form of $C(k_1, k_2)$, and thus, the geometrical parameters obtained with (5) could be misleading. For instance, in the Bjorken Inside-Outside Cascade picture there is a strong correlation between longitudinal coordinate z and longitudinal momentum component p_z[20,21,22,23], which drastically alters the form of the correlation function. In such cases the analysis of correlation functions necessarily becomes model dependent.

In the case of pion interferometry of nuclear collisions at 200 AGeV there are several dynamical effects which may be important to take into account, even if we restrict our attention to completely chaotic sources for which $\lambda = 1$. The rapidity distribution is clearly not uniform and has a width[6] $Y_c \approx 1.4$ for π^-. Also a large fraction of the final π^- could arise from the decay of long lived resonances such as $\omega, K^*, \eta, \cdots$ [12]. In coordinate space, the finite nuclear thickness together with resonance effects can lead to a large spread ($\Delta\tau$) of freeze-out proper times and to a wide distribution of transverse decoupling radii. In phase-space, the imperfect correlation between the space-time and momentum rapidity variables, defined by $\eta = \frac{1}{2}\log((t+z)/(t-z))$ and $y = \frac{1}{2}\log((E+p_z)/(E-p_z))$, should be taken into account. Other correlations, e.g., between the transverse coordinate ($\mathbf{x}_\perp$) and the transverse momentum component ($\mathbf{p}_\perp$), may have to be considered if hydrodynamic flow occurs.

To incorporate these many effects, a Monte Carlo program was developed[7], based on the covariant current ensemble formalism[21]. In this formalism an ensemble of pion source currents is specified as $\{j_a(x) = j_0(u_a^\mu(x-x_a)_\mu)\}$, where u_a^μ is the four-velocity, x_a is the space-time origin of current element a, and $j_0(x)$ specifies each current element in its rest frame. The Fourier transform of the total source current is then

$$j(k) = \sum_a j_0(u_a k) e^{ikx_a} e^{i\phi_a} \quad . \tag{8}$$

where the factors $e^{i\phi_a}$ are random phases in the case of completely chaotic sources. The *double*-pion inclusive distribution function is then given by

$$P_2(k_1, k_2) = \langle | j(k_1) |^2 | j(k_2) |^2 \rangle \quad . \tag{9}$$

which involves the ensemble average over all the space-time coordinates x_a, four velocities u_a, and random phases ϕ_a. In terms of the pion "freeze-out" distribution

$$D(x,p) = \langle \delta^4(x - x_a)\delta^4(p - p_a) \rangle \ , \tag{10}$$

where $p_a^\mu = m u_a^\mu$, the *single* and *double* pion inclusive distribution functions can be written as[21]

$$P_1(k) = G(k,k) \ , \ P_2(k_1, k_2) = G(k_1, k_1)G(k_2, k_2) + \mid G(k_1, k_2) \mid^2 \ . \tag{11}$$

where the complex amplitude $G(k_1, k_2)$ is given by the convolution of the freeze-out distribution and two current elements that characterize the production dynamics,

$$\begin{aligned} G(k_1,k_2) &= \int d^4p D(q,p) j_0^*(pk_1/m) \\ j_0(pk_2/m) &= \langle e^{iqx_a} j_0^*(p_a k_1/m) j_0(p_a k_2/m) \rangle \ , \end{aligned} \tag{12}$$

with $q^\mu = k_1^\mu - k_2^\mu$.

To make specific calculations, the covariant pseudo-thermal parametrization for the current elements[21] was adopted, in which

$$j_0(pk/m) = \exp(-p^\mu k_\mu/(2mT))$$

is characterized by a effective "temperature" T. In this case the amplitude assumes the particularly simple form $G(k_1, k_2) = \langle \exp(iq_\mu x_a^\mu - K_\mu p_a^\mu/(mT)) \rangle$ where $K^\mu = \frac{1}{2}(k_1^\mu + k_2^\mu)$.

To include effects of long lived resonances in the semi classical approximation, we note that the pion freeze-out coordinates, $(t_a, \mathbf{x}_a)$, are related to the resonance production coordinates, $(t_r, \mathbf{x}_r)$, by $\mathbf{x}_a = \mathbf{x}_r + \mathbf{v}_r(t_d - t_r) \ , \ t_a = t_d$, where $\mathbf{v}_r$ is the resonance velocity and t_d is the time of its decay. Averaging over the proper decay time of the resonance, of width Γ_r, leads to

$$G(k_1,k_2) \approx \langle \sum_r f_r \left(1 - \frac{iqp_r}{m_r\Gamma_r}\right)^{-1} \exp(iqx_r - Kp_r/(m_r T_r)) \rangle \ , \tag{13}$$

where f_r is the fraction of the final π^-'s arising from the decay of a resonance of type r, and T_r is an effective temperature that characterizes the decay distribution of that resonance. The ensemble average in (13) is evaluated by Monte Carlo sampling from the following non-ideal freeze out distribution:

$$D(x,p) \propto \tau e^{-\tau^2/\tau_f^2} e^{-(\eta-y)^2/2\Delta\eta^2} e^{-(y-y^*)^2/2Y_0^2} e^{-r_\perp^2/R_\perp^2} \ . \tag{14}$$

To compare with data on the transverse projected correlation function[6], $\langle C(q_\perp)\rangle$,the six dimensional integrals over external momenta (k_1, k_2) were calculated in both numerator and denominator by importance sampling.

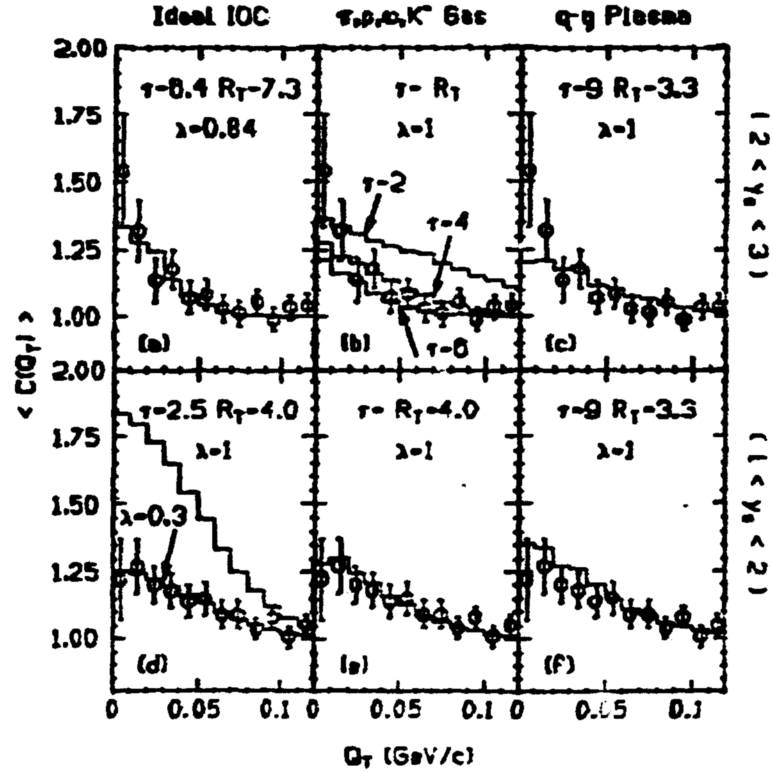

FIGURE 3:

Analysis of the transverse projected $\pi^-\pi^-$ correlation data of NA35[6]. The histograms in parts (a,d) are calculated assuming an ideal inside-outside cascade (IOC) source with parameters $(\tau \equiv \tau_f, R_T \equiv R_{\perp f})$ taken from [6]. In parts (b,e) a non-ideal resonance gas source is considered with parameters, $\tau \sim R_T \sim 4$ fm, as suggested by the ATTILA version of the LUND Fritiof model[25]. Parts (c,f) correspond to the quark-gluon plasma model of [24]. Parts (a-c) refer to the central rapidity region, $2 < y_\pi < 3$, and parts (d-f) refer to the region $1 < y_\pi < 2$.

A good test of the numerical method of Ref.[7] is provided by reproducing the fitted curves in Ref.[6], which follow assuming the ideal inside-outside cascade distribution. Fig 3a and 3d show the calculations employing the reported parameters[6], $\tau_f = 6.4$ fm/c, $R_\perp = 7.3$ fm and $\lambda = 0.84$ for π^- in the rapidity interval $2 < y < 3$ and $\tau_f = 2.5$ fm/c, $R_\perp = 4.0$ fm, and $\lambda = 0.30$

in the interval $1 < y < 2$, which in fact provide a good fit to the data (note that the data have been corrected for Coulomb final state interactions).

Next, Figs 3b,3c show the calculated curves for the case of non-ideal hadron resonance dynamics discussed above. For these calculations the values of $\Delta\eta = 0.8$ and considered $\tau_f = R_\perp = 2, 4, 6$ fm were chosen. An additional Monte Carlo hadronic cascade calculation was performed taking as input the output of the LUND fragmentation model[25] and found that with a $\sigma = 20$ mb, the true freeze-out distribution is roughly characterized by $\tau_f \sim R_\perp \sim 4$ fm for this reaction. In both Fig. 1b and 1e, the chaoticity parameter is fixed to $\lambda = 1$ as appropriate for completely chaotic sources. It is clear that the present data are consistent with the freeze-out distribution expected on the basis of a resonance gas model for the nuclear dynamics.

Next, Figs. 3c,3f show the remarkable result that the quark-gluon plasma freeze-out distribution is also consistent with the data. The reason is that the long lifetime of the plasma source leads to the same effect in this case as the inclusion of long lived hadronic resonances in Figs. 3b,3e. While it would be difficult to justify ruling out any of the three models from the present data, the "exotic" parameters obtained with the ideal inside-outside cascade model[21] are the least compelling, since it would be truly remarkable if the degree of coherence in such violent nuclear collisions were not negligible.

The coincidental agreement with data of the two opposite scenarios, such as a resonance gas and a quark gluon plasma mentioned above stresses the necessity of finding other means to more clearly differentiate among dynamical models.

Complementary studies of the constructive interference pattern of bosons other than pions could help to enhance the distinction. A suggestion was made[28] to compare pion and kaon interferometry. The main motivation for this is that an entirely different set of hadronic resonances decay to kaons than to pions, thus leading to a completely different freeze out geometry. In particular, in the resonance gas model, the freeze out proper time for kaons is expected to be much smaller than for pions due to the fact that long lived ω, η and η' do not contribute to the first ones. According to the ATTILA/LUND fragmentation model[25], approximately one half of the kaons are produced by direct string decay and the other half, by the decay of K^*.

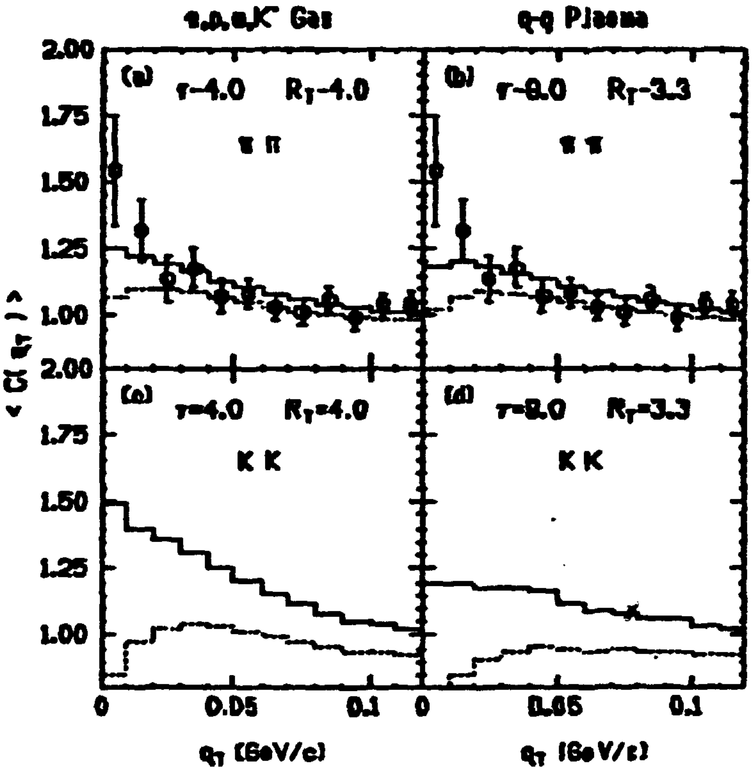

FIGURE 4:

In (a),(b) and (c),(d) we have, respectively, pion and kaon projected correlation functions versus transverse momentum difference q_T, for $q_L \leq 0.10 GeV/c$, in central rapidity region. (a),(c) correspond to predictions based on the ATTILA/ LUND model[25]. Parts (b,d) correspond to the plasma hydrodynamic model of Ref.[24]. The pion data are from Ref.[6].

On the other hand, in the plasma model of Ref.[24], the freeze out geometry of all hadrons is expected to be about the same, and controlled mainly by the slowness of the deflagration process. The resulting transverse momentum projected correlation functions are shown in Fig.4 for the central rapidity region. A finite cut on the longitudinal momentum $q_L \leq 0.1 GeV/c$ was taken. The solid (dashed) curves indicate correlation functions without (with) Coulomb-Gamow corrections. As already mentioned in the beginning. Figs. 4 (a) and (b) reveal that, for O + Au at 200 A GeV and the corresponding kinematical cuts, neither the resonance or the plasma model could be ruled out by the experimental data. An entirely different behavior is expected in the case of kaons: the calculated curves show a much more significant difference between these two models. The major disadvantage of kaon interferometry is the need for vastly higher statistics for kaons $[(\pi/K)^2 \approx 100]$. In addition, final state Coulomb interactions, which is bigger for kaons than for pions due to the increase of mass. However, Coulomb corrections are well understood.

In the "non-ideal" resonance gas model discussed here, resonances are the major responsible for the distorting effects caused on the correlation function (see Ref.[29]). The resonance fraction is in fact of interest in its own right. Final state dynamical effects could change the resonance abundance predicted by string models. Concerning this point, we show in Fig. 5 results for pion interferometry with the experimental cuts corresponding to the AGS run on $Si + Au$ at 14.6 A GeV. The continuous histograms correspond to inclusion of resonances and the dashed ones to interferometry without resonances[30]. Fig. 5 also shows the histograms corresponding to two q_L bins: the lower ones are for the bigger interval $q_L \leq 0.10 GeV/c$ and the higher ones, are related to the projection of the first bin only, for which $q_L \leq 0.01 GeV/c$ (as in Ref.[31]).

The comparison of the results for these two binnings, nicely illustrates the striking influence of finite q_L on lowering the correlation function intercept: even without the inclusion of resonances, finite binning causes a dramatic decrease of the apparent intercept, a fact emphasized in [7,22,28]. In Fig.5 a clear sensitivity to resonances can be seen. A first comparison with the preliminary experimental data of Ref.[31] appears to suggest the near absence of resonances in AGS range as compared to SPS energies. It will be of considerable interest to follow up this possibility. The point here is that in the absence of direct resonance measurements, pion interferometry provides at least an indirect probe of their abundance.

Finally, we show here the result of the study of the sensitivity of different reactions to resonances, analyzing in this way the trends of the freeze out geometry as a function of the atomic number.

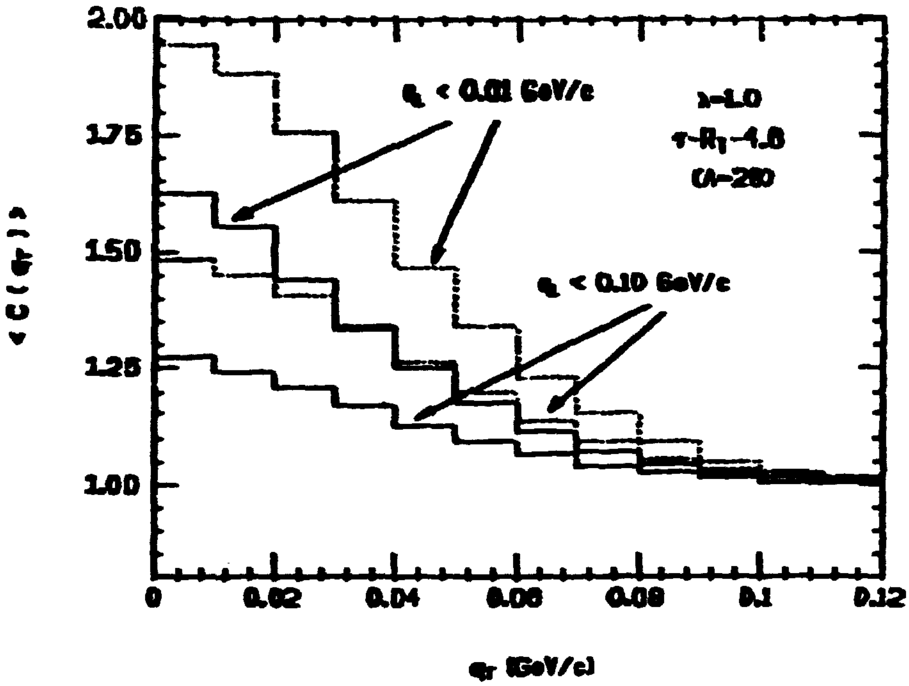

FIGURE 5:

We can see here the predictions for $< C(q_T) >$ versus q_T at AGS energies. Continuous (dashed) histograms correspond to resonances (not) included in the analysis. For the two lower results, $q_L < 0.10 GeV/c$ and for the higher ones, $q_L < 0.01 GeV/c$.

In Fig. 6 we show a preliminary calculation for increasing A (= 32, 64, 128 and 197) where we respect the same kinematical cuts as in Ref.[6]. The radius was estimated as $R_{Tf} \approx R_0 A^{1/3} + 1(fm)$, the proper time $\tau_f \approx R_{Tf}$ with inclusion of all the non ideal dynamical and geometrical parameters. The extra factor of $1 fm$ is included to take into account possible final state cascading. This extra factor is, however, highly dependent on the hadronic density and a possible change with A is under investigation.

In Fig. 6, comparison is made between resonances included and in their absence for each value of A. We notice that the differences between these two cases decrease with increasing A. The trends of the behavior of the correlation function $C(q_T)$ versus q_T can be seen: it becomes narrower and the intercept drops when A is enhanced. These results are expected since, as A increases, the correlation effect is limited to an ever decreasing phase-space. The dropping intercept again reflects the finite q_L binning considered, as already discussed above.

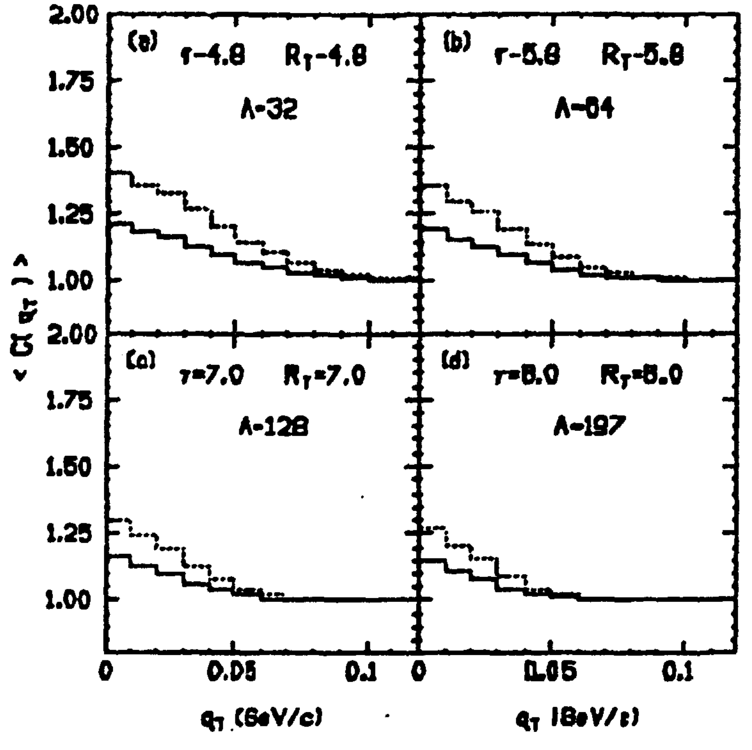

FIGURE 6:

The correlation function $< C(q_T) >$ versus q_T is shown for increasing atomic number, considering $q_L < 0.10 GeV/c$. The discussed non-ideal features are incorporated into the analysis. Continuous histograms correspond to resonances included and dashed ones to no resonances.

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