Title
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Permalink
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Journal
Physics Letters B, 286

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Publication Date
1992-04-01

LBL-32051
UC-413
Preprint

Lawrence Berkeley Laboratory
UNIVERSITY OF CALIFORNIA

Submitted to Physics Letters B

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April 1992

NUCLEAR SCIENCE DIVISION

Prepared for the U.S. Department of Energy under Contract Number DE-AC03-76SF00098

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LBL-32051

K_s Pictures of Strangeness Distillation*

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April 12, 1992

Abstract:
$K_s - K_s$ interferometry is shown to contain a novel interference term that is sensitive to possible strangeness distillation phenomena in high energy nuclear reactions. The absence of Coulomb final state interaction, the reduced resonance distortions, and the above novel effect make it an especially attractive complement to conventional charged boson ($\pi^-\pi^-$ or K^+K^+) interferometry.

Pion and kaon interferometry are used to probe the space-time geometry of hadronic and nuclear reactions [1]-[11]. In the ideal case, symmetrization of identical boson wavefunctions induces a correlation function, $C_2(\mathbf{k}_1, \mathbf{k}_2) = 1 + |\rho(k_1 - k_2)|^2$, that is related to the Fourier transform of the *decoupling* space-time density, $\rho(x)$, of the produced particles. However, in most cases, dynamical effects lead to strong correlations between spatial and momentum coordinates which distort the interference pattern and obscure the space-time interpretation of $C_2(\mathbf{k}_1, \mathbf{k}_2)$[7]. In addition, distortions due to resonance production[5] and final state interactions[3, 12] must be taken into account before any geometric information can be deduced. In Ref.[7] a general semi-classical formulation of the problem was derived. In this note we adapt that formalism to derive a new twist on this old topic in connection with $K_s - K_s$ interferometry.

The motivation for concentrating on K_s stems from the following considerations: (1) Coulomb final state interaction distortions are absent, (2) the distortions due to resonance formation are less severe than for pions[6], (3) it can probe possible novel aspects of the strangeness production dynamics in nuclear collisions [13, 14, 15], and (4) it appears to be experimentally feasible with a resolution not limited by the ionization width of overlapping charged particle tracks [16]. The main physics motivation

[0]*This work was supported by the Director, Office of Energy Research, Division of Nuclear Physics of the Office of High Energy and Nuclear Physics of the U.S. Department of Energy under Contract No. DE-AC03-76SF00098.

suggested in [13, 14, 15], is that a system at finite baryon densities may under certain conditions distill strange quarks by evaporating K^+ and K_0 preferentially to K^- and $\overline{K}_0$. In a quark-gluon plasma strangeness distillation could occur if the chemical potential of light quarks exceeds the strange quark mass. In a baryon rich hadronic gas, the difference between $\sigma(K^+p) \sim 10$ mb and $\sigma(K^-p) \sim 50$ mb for $p_{rel} \lesssim 1$ GeV can cause a slower diffusion of the negative strangeness in high baryon density matter[15]. In either case, one expects a relative time delay between the K_0 and $\overline{K}_0$ emission. In Refs.[14, 13] the emphasis was on the difference between $\Lambda\Lambda$ and $\pi\pi$ interferometry to study this effect. K_s interferometry was only briefly considered, and as we show below, an interesting interference effect arising from the weak mixing of strangeness and anti-strangeness was missed.

In [6] the advantage of kaon versus pion interferometry was demonstrated in connection with reduced resonance distortions. However, in [6] only charged kaon pairs were considered, for which unfortunately the Coulomb distortions are more severe than for pions at the same momentum difference. Recall that the multiplicative Gamov factor[3], $\chi(q) = (q_c/q)/(e^{q_c/q} - 1)$ where $q_c = 2\pi\alpha m$ drills a hole in the correlation function of width, $q_c \sim 23MeV$, which is 3.5 times as broad for kaons than for pions. This leads to $C < 1$ for charged kaon pairs for $q \lesssim q_c$ as shown in Fig.1 of [6]. If the decoupling time or radius exceeds 10 fm, then all the sought after interference phenomena must be extracted from the Coulomb hole! For three or more charged particle interferometry[17] the Coulomb distortion becomes more severe and theoretically poorly understood except in the perturbative domain at high q, where it is negligible. On the other hand, the main practical advantage of charged mesons is that they are usually easier to detect and at least the two particle Coulomb correction is well understood. In any case, it is clearly useful to measure the Coulomb free K_s correlations, especially because of its novel sensitivity to strangeness distillation phenomena[13, 14, 15] as demonstrate below. Its main advantage over $\Lambda\Lambda$ interferometry is that the later is severely distorted by as yet poorly understood strong final state interactions in addition to resonance distortions[14].

The Wigner density formalism[18] as applied in ref.[7] provides a convenient starting point for this problem. The inclusive number distribution, $P(\alpha)$, of any asymptotic multiparticle state, $|\alpha\rangle$, can be calculated formally from the exact density matrix, $\rho(t) = |\psi(t)\rangle\langle\psi(t)|$ as

$$\begin{aligned} P(\alpha) &= \lim_{t\to\infty} Tr\rho_\alpha \rho(t) = \int dt Tr\rho_\alpha \frac{d}{dt}\rho(t) \\ &= -i\int dt Tr\rho_\alpha [H_1(t), \rho(t)] \equiv -i\int dt Tr\rho_\alpha \left(\frac{d}{dt}\right)_I \rho(t) \ , \end{aligned} \tag{1}$$

where the Hamiltonian, $H = H_0 + H_1$, was split into a part, H_0, whose eigenstates include, $|\alpha\rangle$, and an interaction part, H_1, responsible for the production of the particles and high momentum transfer interactions. In the case of K_s, H_0 must include weak interactions which mix the K_0 and $\overline{K}_0$ strangeness eigenstates states of the strong interaction such that up to tiny CP violation $|K_s\rangle = (|K_0\rangle + |\overline{K}_0\rangle)/\sqrt{2}$.

Evaluating the trace in the Wigner representation, the 2-particle inclusive distribution is

$$P(\mathbf{k}_1, \mathbf{k}_2) = \int dt \int d\phi_1 d\phi_2 W_{\mathbf{k}_1,\mathbf{k}_2}(\phi_1, \phi_2) \left(\frac{d}{dt}\right)_I f_2(\phi_1, \phi_2, t) \ , \tag{2}$$

where $\phi_i = (\mathbf{x}_i, \mathbf{p}_i)$ are six dimensional phase space coordinates with integration measure $d\phi_i = d^3\mathbf{x}_i d^3\mathbf{p}_i (2\pi)^{-3}$ in standard ($\hbar = c = 1$) units. It is clear from (2) that only the rate of change of the two particle density, f_2, on the strong interaction time scale due to H_I influences the asymptotic inclusive yield. That rate of change can be approximated using an appropriate transport theory[18]. The Wigner representation of the asymptotic density matrix in the CP eigenbasis where, $\rho_\alpha = |K_s(\mathbf{k}_1)K_s(\mathbf{k}_2)\rangle\langle K_s(\mathbf{k}_1)K_s(\mathbf{k}_2)|$, is

$$W_{\mathbf{k}_1,\mathbf{k}_2}(\mathbf{x}_1, \mathbf{p}_1, \mathbf{x}_2, \mathbf{p}_2) = \delta^3(\mathbf{p}_1 - \mathbf{k}_1)\delta^3(\mathbf{p}_2 - \mathbf{k}_2) + e^{-i\mathbf{q}(\mathbf{x}_1 - \mathbf{x}_2)}\delta^3(\mathbf{p}_1 - \mathbf{K})\delta^3(\mathbf{p}_2 - \mathbf{K}) \ , \tag{3}$$

where $\mathbf{q} = \mathbf{k}_1 - \mathbf{k}_2$ and $\mathbf{K} = (\mathbf{k}_1 + \mathbf{k}_2)/2$. The interference term is due to Bose symmetry.

The new twist about neutral kaon interferometry is that the rate of change of the two kaon phase space density, f_2, is diagonal in the K_0, $\overline{K}_0$ strangeness basis during the strong interaction reaction time. Thus, in the asymptotic CP basis

$$\langle K_s K_s || f_2 || K_s K_s \rangle = \frac{1}{4}(f_{++} + f_{+-} + f_{-+} + f_{--}) \ , \tag{4}$$

where $\pm$ refers to the strangeness quantum number. In Ref.[14] the cross terms, $f_{\pm\mp}$, leading to the novel interference effect discussed below were neglected.

To proceed further, we solve the transport equation for the f_{ij} as in [7] and integrate (2) over time. In the quasi-equilibrium limit where two-particle *dynamical* correlations can be neglected and the contributions from intermediate scattering times can be neglected, the correlation function depends[5, 7] only on the single particle *decoupling* distributions $D_+(x,p)$ and $D_-(x,p)$ of the K_0 and $\overline{K}_0$ respectively:

$$C_2(\mathbf{k}_1, \mathbf{k}_2) = 1 + \frac{|D(q,K)|^2}{D(0,k_1)D(0,k_2)} \ , \tag{5}$$

where K and q are now the average and difference of the four momenta, k_1 and k_2, and the invariant distribution amplitude can be expressed as

$$\begin{aligned} D(q,K) &= \int d^4x d^4p \, e^{iqx} e^{-Kp/\Delta^2} (D_+(x,p) + D_-(x,p)) \\ &\equiv D_+(q,K) + D_-(q,K) \ . \end{aligned} \tag{6}$$

These decoupling distributions are the solutions to the underlying transport equations. For example, in a classical cascade transport model

$$D_\pm(x,p) = \langle \delta^4(x - x_f)\delta^4(p - p_f) \rangle_\pm \ , \tag{7}$$

in terms of the space time point, x_f, of the final high momentum transfer scattering event in the comoving matter that leaves the K_0 or $\overline{K}_0$ with a final momentum,

p_f. Here $\langle ... \rangle$ denotes the ensemble average over all possible trajectories. In a hydrodynamic model, these distributions are usually fixed on an assumed decoupling hypersurface, e.g. $T(x_f) = T_f$.

The Δ appearing in (6) is a momentum dispersion and arises when minimal wavepackets are used to smear over the classical phase space coordinates, (x_f, p_f), to take the uncertainty principle into account[7]. As emphasized in [7], $D(q,K)$ becomes independent of Δ only in the true semi-classical limit. However, very strong correlations expected between coordinates and momenta in high energy reactions (the inside-outside cascade) can lead to a dependence of D on Δ. Fortunately, that dependence is strongly constrained by the single inclusive distribution[5]. The often used[11, 14] ansatz[2], obtained by setting $\Delta \to \infty$ and replacing $|D(q,K)|^2$ by $ReD(q,k_1)D(-q,k_2)$ is valid only in the strict semi-classical limit, where $\hbar\partial_x\partial_p \log D(x,p) \ll 1$, and may significantly underestimate the rate of change of the correlation function with respect to the rapidity difference[7].

To proceed, note that $D_\pm(q,K)$ can be identified with the invariant distribution amplitude entering $K^\pm K^\pm$ interferometry if the K^+ and K_0 decoupling densities are assumed to be identical and that the K^- and $\overline{K}_0$ decoupling densities are assumed to be identical. In this case, the single inclusive distributions are

$$\begin{aligned} \omega d^3N(K^+)/dk^3 &= \omega d^3N(K_0)/dk^3 \propto D_+(0,k) \ , \\ \omega d^3N(K^-)/dk^3 &= \omega d^3N(\overline{K}_0)/dk^3 \propto D_-(0,k) \ . \end{aligned} \tag{8}$$

Defining, the reduced (Gamov corrected) identical kaon correlation function

$$R_a(q,K) = C_a(k_1,k_2)/\chi_a(q) - 1 \ , \tag{9}$$

where $a = s,+,-$ for K_s ,K^+, and K^-, resp., and $\chi_s = 1$, $\chi_\pm = \chi$, we see that

$$\begin{aligned} R_\pm(q,K) &= |D_\pm(q,K)|^2/(D_\pm(0,k_1)D_\pm(0,k_2)) \ , \\ R_s(q,K) &= f(k_1)f(k_2)R_+(q,K) \ + \ (1-f(k_1))(1-f(k_2))R_-(q,K) \\ &\quad +2f(k_1)(1-f(k_2))I_{+-}(q,K) \ . \end{aligned} \tag{10}$$

The local K^+ fraction is the ratio of single inclusive cross sections given by

$$f(k) = D_+(0,k)/D(0,k) = K^+/(K^+ + K^-) = K_0/(K_0 + \overline{K}_0) \ , \tag{11}$$

and the new interference term is

$$I_{+-}(q,K) = Re(D_+(q,K)D_-(-q,K))/(D_+(0,k_1)D_-(0,k_2)) \ . \tag{12}$$

Note that the intercept of the reduced K_s correlation function at $q=0$ reaches the ideal Bose limit, $R_s(0,K) = 1$, unlike in the ansatz used in [14]. Also, at zero baryon density, $D_+ = D_-$, and all the reduced correlation functions are identical. However, for non-vanishing baryon density the difference is amplified via the new term. Thus, unlike in [14], where for $f = 1-\epsilon$, corrections to $R_s = R_+$ appear to $O(\epsilon^2)$, the derivation above shows that they are $O(2\epsilon)$. (For $K_L K_L$ interferometry only the sign of I_{+-} must be reversed.)

To illustrate the above formula, we consider an ideal inside-outside cascade decoupling phase space density[7],

$$D_\pm(x,p) \propto N_\pm \delta(\tau(t,z)-\tau_\pm)\delta(\eta - y)\delta(E-E_{\rm p})\delta^2(\mathbf{p}_\perp)e^{-x_\perp^2/r_\pm^2} \ , \tag{13}$$

where $\tau(t,z) = (t^2-z^2)^{\frac{1}{2}} = \tau_\pm$ is the decoupling *proper time* for $K^\pm$, resp. and the space-time and momentum rapidity variables,

$$\eta = \tfrac{1}{2}\log((t+z)/(t-z)) \ , \ y = \tfrac{1}{2}\log((E+p_z)/(E-p_z)) \ , \tag{14}$$

are assumed to be perfectly correlated. Above we also allow for the rms transverse decoupling radii, $r_\pm$, to differ. The possibility that the $K^\pm$ transverse momentum distributions differ[15] can be taken into account by setting $\Delta^2 = mT_\pm$ so that

$$\omega d^3N(K^\pm)/dk^3 \propto D_\pm(0,k) \propto N_\pm e^{-m_\perp/T_\pm}/\sqrt{m_\perp} \ . \tag{15}$$

The correlation function can be expressed analytically in terms of Bessel functions[7], but it is most instructive to look at the kinematic region for which the transverse masses are equal , $m_{\perp 1} = m_{\perp 2} = m_\perp$ and large compared to $T_\pm$, the rapidity difference $y_1 - y_2$, is small, but otherwise allow the transverse momentum difference, $\mathbf{q}_\perp = \mathbf{k}_{\perp 1} - \mathbf{k}_{\perp 2}$ to vary. In this limit

$$D_\pm(q,K) \approx N_\pm e^{-q_\perp^2 r_\pm^2/4}(1 - \tfrac{1}{2}\tau_\pm^2 T_\pm m_\perp (y_1-y_2)^2) \ , \tag{16}$$

and

$$f = f(k_1) = f(k_2) = N_+ e^{-m_\perp/T_+}/(N_+ e^{-m_\perp/T_+} + N_- e^{-m_\perp/T_-}) \ . \tag{17}$$

These kinematics simplify tremendously the reduced correlators giving

$$\begin{aligned} R_\pm(q,K) &\approx e^{-q_\perp^2 r_\pm^2/2}(1 - \tau_\pm^2 T_\pm m_\perp (y_1-y_2)^2) \ , \\ R_s(q,K) &\approx f^2 R_+(q,K) + (1-f)^2 R_-(q,K) \\ &\quad +2f(1-f)e^{-q_\perp^2(r_+^2+r_-^2)/4}(1 - \tfrac{1}{2}(\tau_+^2 T_+ + \tau_-^2 T_-)m_\perp(y_1-y_2)^2) \end{aligned} \tag{18}$$

For high baryon densities, we expect $f = 1-\epsilon$, where $\epsilon \approx K^-(k_\perp)/K^+(k_\perp)$. To first order in ϵ

$$R_s(q,K) \approx (1-2\epsilon)R_+(q,K) + 2\epsilon e^{-q_\perp^2(r_+^2+r_-^2)/4}(1 - \tfrac{1}{2}(\tau_+^2 T_+ + \tau_-^2 T_-)m_\perp(y_1-y_2)^2) \ . \tag{19}$$

This shows that the K_s correlation function may have a two component structure if strangeness distillation occurs. One component is just R_+ suppressed by the factor $1-2\epsilon$, and the other depends on the decoupling parameters of both strangeness and anti-strangeness. The second component has a magnitude 2ϵ and is localized to

$$\begin{aligned} \Delta q_\perp &\sim 2/(r_+^2 + r_-^2)^{1/2} \\ \Delta y &\sim (2/(\tau_+^2 T_+ + \tau_-^2 T_-)m_\perp)^{1/2} \ . \end{aligned} \tag{20}$$

Since $T_\pm$ is fixed by the measured single inclusive transverse momentum spectrum, separate measurements of the $q_\perp$ and $y_1 - y_2$ dependence of R_+ and R_s can thus in principle determine the four unknown decoupling scales $r_\pm$ and $\tau_\pm$. Since we expect $r_+ < r_-$ and $\tau_+ < \tau_-$, the second component should to be narrower than the dominant K^+ component. Note however, that the difference between $\tau_\pm$ can be partially offset by the difference between $T_\pm$ which may also be associated with strangeness distillation[15].

The interference pattern is of course complicated by K^* resonance formation[6, 19]. Typically a large fraction, $f^* \approx 0.5$ of the final kaons emerge from decays of K^*. Fortunately, other resonance channels are negligible. Taking only K^* into account modifies the amplitude (6) as[5, 7]

$$D_\pm(q,K) = \int d^4p \left\{ (1-f^*)e^{-Kp/\Delta^2} D^d_\pm(q,p) + f^* \frac{e^{-Kp/\Delta^{*2}} D^*_\pm(q,p))}{(1 - iqp/m^*\Gamma^*)} \right\} \quad , \qquad (21)$$

where $D^d_\pm(q,p)$ and $D^*_\pm(q,p)$ are the spatial Fourier transforms of the direct and K^* decoupling distributions, and $\Delta^{*2} = \Delta^2 + p_0^2$ with $p_0 \approx 0.29$ GeV is the momentum kick of the K_s due to the K^* decay ($m^* = 894, \Gamma = 50$ MeV). Expanding (21) to second order in q shows that resonance production increases the apparent lifetime and size of the source.

Resonance production obviously influences the form of the K^+K^+ correlation as well. Numerically, it was found[6] that K^* production makes the correlation more cusp-like near $q = 0$ then a naive Gaussian parameterization would suggest. The important point is that this resonance effect already shows up in the K^+K^+ correlation. Therefore with an independent measurement of K^+K^+ correlations, the uncertainties associated with f^* and $D^*_+(q,p)$ can be considerably reduced.

In practice, hadron interferometry is most useful as a tool to test specific transport models[13, 15] that predict different $f^*, D^d_\pm(x,p), D^*_\pm(x,p)$. The first test of any such model must be the single kaon inclusive spectra and the K^+K^+ correlation function. If successful, the novel extra interference in K_s can then be used to test the predicted kaon distillation phenomena when K^-K^- is impractical. Even when K^-K^- is practical, the novel quantum mechanical interference arising due to weak mixing of strange eigenstates combined with Bose symmetrization would be interesting to verify.

The most suitable energy range for such studies is $E_{lab} \sim 10 - 15$ AGeV (at the AGS) where maximal baryon densities are likely to be achieved in $Au + Au$ collisions and K^+K^+ interferometry has already been performed[20]. At these energies, $\epsilon = K^-/K^+ \sim 0.2$ so that pure K^-K^- interferometry is much more difficult than K^+K^+. In this case K_sK_s interferometry may be the only practical way to get a glimpse of the anti-strangeness decoupling distribution. Such measurements may also be of interest at SPS and RHIC energies ($\sqrt{s} \sim 20, 200$ AGeV) as an independent check on the K^+K^+ correlation without the Coulomb hole[16] and to search for kaon distillation if unexpectedly high baryon densities turn out to be generated[21].

I am grateful to Declan Keane for discussions and suggesting this problem. Discussions with S.S. Padula and W. Zajc are also gratefully acknowledged.

References

[1] G. Goldhaber, et al., Phys. Rev. 120 (1960) 300.

[2] F. B. Yano, S.E. Koonin, Phys. Lett. 78B (1978) 556.

[3] M. Gyulassy, S. K. Kauffmann, L. W. Wilson, Phys. Rev. C20 (1979) 2267.

[4] S. Pratt, Phys. Rev. Lett. 53 (1984) 1219.

[5] S.S. Padula, M.Gyulassy, Phys. Lett. B217 (1989) 181

[6] M. Gyulassy, S.S. Padula, Phys. Rev. C41 (1989) R21.

[7] S.S. Padula, M. Gyulassy, and S. Gavin, Nucl. Phys. B329 (1990) 357; Nucl. Phys. B339 (1990) 378.

[8] S. Pratt, Phys. Rev. D33 (1986) 1314.

[9] S. Pratt, Phys. Rev. D33 (1986) 72.

[10] G. Bertsch, M. Gong, M. Tohyama, Phys. Rev. C37 (1988) 1896; G. Bertsch, Nucl. Phys. A498 (1989) 173c.

[11] T. Csorgo, et al,Phys. Lett. B241 (1990) 301; S. Pratt, et al, Phys. Re. C42 (1990) 2646.

[12] M. G. Bowler, Z. Phys. C39 (1988) 81.

[13] C. Greiner, P. Koch, H. Stocker, Phys. Rev. Lett. 58 (1987) 1825; Phys.Rev.D38 (1988) 2797.

[14] C. Greiner , B. Muller Phys. Lett. B219 (1989) 199.

[15] U. Heinz , K.S. Lee, M.J. Rhoades-Brown Mod. Phys. Lett. A2 (1987) 153; Phys. Rev. Lett. 58 (1987) 2292.

[16] D. Keane, private communication.

[17] Y. M. Liu, et at, Phys. Rev. C34 (1986) 1667.

[18] E. A. Remler et al., Ann. Phys. 91 (1975) 295; 95 (1975) 455; 136 (1981) 293.

[19] A.M. Cooper et al., Nucl. Phys. **B139**, 45 (1978).

[20] R.J. Morse, MIT thesis 1990; W. Zajc, Proc. QM91.

[21] A. von Keitz, et al, Phys.Lett.B263 (1991) 353.

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