

Effects of collective potentials on pion spectra in relativistic heavy-ion collisions

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Abstract: The effect of collective potentials on pion spectra in ultrarelativistic heavy-ion collisions is investigated. We find the effect of these potential to be very small, too small to explain the observed enhancement at low transverse momenta.

1. Introduction

The enhancement at low transverse momentum found in negative-particle and neutral-pion spectra from ultrarelativistic heavy-ion collisions¹⁻³⁾ has recently received considerable interest in the literature. One explanation for instance^{4,5)} has interpreted the enhancement as a collective-flow effect, but more detailed considerations concerning the freeze-out surface result in too small an effect⁶⁾. The decay of excited baryons also gives rise to additional soft pions⁷⁻⁹⁾. However, at CERN energies there are simply not enough baryons present in the central region in order to account for the observed enhancement⁸⁾. Furthermore the combined spectrum of thermal plus decay pions still flattens out at very low p_t contrary to the data.

Kataja and Ruuskanen¹⁰⁾ have shown that the measured spectrum can be fitted by assuming that the pions are strongly out of equilibrium. Although one may give some qualitative arguments¹¹⁾ why an excess of pions should build up during the expansion, currently there is no quantitative understanding on how a chemical potential of the size needed to fit the data ($\mu \simeq 130$ MeV) would arise.

In refs.^{12,13)} a kinetic model with Bose statistics in the collision integral has been used to study the build up of the low- p_t enhancement assuming different initial conditions for the expanding source. The authors find that the Bose statistics in the collision integral indeed leads to an enhancement provided the initial pion density

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is sufficiently high. They also point out that corrections of the cross section due to Bose statistics are large and that they limit the effect of the Bose phase-space factors in the collision integral substantially. Using the corrected cross section and allowing for higher resonances in the initial state the authors have to assume a hadronization time as short as $\tau_0 \sim 1$ fm in order to account for the measured enhancement by the hadronic scattering processes only.

Finally the pions may interact collectively with the surrounding hadronic medium. As proposed by Shuryak^{14,15)} these interactions may give rise to a strongly momentum-dependent optical potential which is attractive for low momenta and, therefore, could lead to an enhanced soft component in the pion spectrum.

It is the purpose of this article to study the effect of this latter mechanism in detail. In the first part we will develop the mean-field potential and demonstrate how it can be used in a transport theoretical framework. Then we will study the effect of the mean fields for a static potential. Finally the expansion of the fireball will be taken in account. This will be done using a transport model which combines the propagation of particles in the mean field as well as the collisions among the particles.

2. Mean-field potential

Following Shuryak¹⁴⁾, the most important contribution to the collective potential felt by a pion is the coherent scattering of two pions via the P-wave ρ -resonance; in the S-wave the contributions from different isospin channels cancel each other. We first show that these collective potentials can be derived from an energy functional by differentiating with respect to the distribution function. As a consequence energy conservation will be guaranteed and the collective potentials can be incorporated similar to the nuclear mean field in transport models of heavy-ion collisions^{16,17)}.

For an arbitrary resonance the energy functional has the following form:

$$H = g_{I,J} \int \frac{d^3P}{(2\pi)^3 2\omega(p)} \frac{d^3p'}{(2\pi)^3 2\omega(p')} \sum_i f_i(x, p) \times \sum_m f_m(x, p') \frac{4\pi\sqrt{s}}{q} \frac{(\sqrt{s} - M)\Gamma(q)}{(\sqrt{s} - M)^2 + \frac{1}{4}\Gamma^2(q)}, \quad (1)$$

with q and \sqrt{s} being the c.m. momentum and energy and M and Γ are the mass and the width of the resonance under consideration e.g. the ρ . The phase-space distribution $f_i(x, p)$ contains an index which refers to the internal quantum numbers such as spin and isospin. These quantum numbers are summed over. The degeneracy factor $g_{I,J}$ is given by

$$g_{I,J} = \frac{1}{1 + \delta_{p_1, p_2}} \frac{(2J_{\text{res}} + 1)(2I_{\text{res}} + 1)}{(2J_{p_1} + 1)(2I_{p_1} + 1)(2J_{p_2} + 1)(2I_{p_2} + 1)}, \quad (2)$$

where the indices p_1 and p_2 refer to the particles forming the resonance. For example in the case $\pi + \pi \rightarrow \rho$ we would have $g_{l,j} = \frac{1}{2}$.

The resulting mean-field potential for a particle of given internal quantum number l is then given by

$$U_l(p) = \frac{\delta H}{\delta f_l} = g_{l,j} \frac{1 + \delta_{p_1, p_2}}{2\omega(p)} \int \frac{d^3 p'}{(2\pi)^3 2\omega(p')} \times \sum_m f_m(x, p') \frac{4\pi\sqrt{s}}{q} \frac{(\sqrt{s} - M)\Gamma(q)}{(\sqrt{s} - M)^2 + \frac{1}{4}\Gamma^2(q)}, \quad (3)$$

which in the special case of $\pi + \pi \rightarrow \rho$ leads to the result already obtained by Shuryak. Assuming isospin symmetry we obtain

$$U_\pi(p) = \frac{3}{2\omega(p)} \int \frac{d^3 p'}{(2\pi)^3 2\omega(p')} \bar{f}(x, p') \frac{4\pi\sqrt{s}}{q} \frac{(\sqrt{s} - M_\rho)\Gamma_\rho(q)}{(\sqrt{s} - M_\rho)^2 + \frac{1}{4}\Gamma_\rho^2(q)}, \quad (4)$$

where $\bar{f}(x, p) = \frac{1}{3} \sum_{l=+,-,0} f_l(x, p)$ stands for the isospin-averaged pion phase-space distribution.

In the following we will truncate the explicit energy dependence of the potential U by evaluating eq. (3) on the mass shell only, i.e. $\omega^2 = p^2 + m_\pi^2$, which is equivalent to taking only the first contribution in the Dyson series for the selfenergy. As a consequence the potential depends on the three-momentum of the particle only and, therefore, the pion wave function will not be modified. Hence we can treat the mesons as quasiparticles and a transport theoretical approach is possible*.

For local equilibrium the distribution function can be written as

$$f(x, p) = \rho_0 \exp(-E/T(x)), \quad (5)$$

where $E = \sqrt{p^2 + m^2}$ denotes the energy and $T(x)$ the local temperature. Using the above phase-space distribution (5) we can calculate the mean-field potential $U(p)$ from eqs. (3) or (4).

In fig. 1 the resulting pion collective potential due to $\pi + \pi \rightarrow \rho$ are shown for three different temperatures ($T = 150, 200, 250$ MeV) (full lines). We find the potential becomes deeper with increasing temperature, because at the same time the density of pions in the heat bath increases. The point where the potential changes sign, on the other hand, is essentially unaffected by the temperature.

Since we are interested in a more general discussion of mean-field effects in this article, in fig. 2 we show the dependence of the collective pion potential on the resonance mass. These potentials have been obtained for the same temperature $T = 200$ MeV but with different resonance mass $M_{\text{res}} = 500, 770, 1100$ MeV. For the width we have taken a p-wave parameterization with a value of $\Gamma_0 = 150$ MeV on resonance and the degeneracy factor was chosen to be the same as in $\pi + \pi \rightarrow \rho$. As we would have expected from eq. (3) the potentials change sign at a momentum close to the resonance mass.

* Note that the formalism described here may be used as well for pions interacting with nucleons to study in-medium effects on pions in lower-energy collisions.

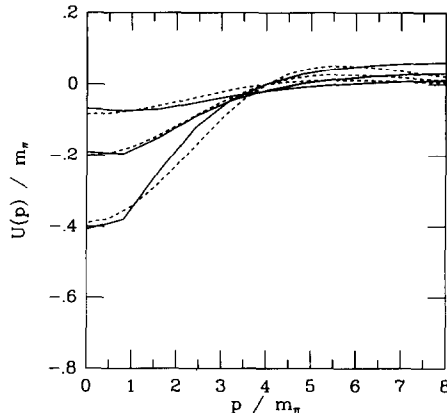


Fig. 1. Mean-field potential for $M_{\text{res}} = 770$ MeV for different temperatures ($T = 150, 200, 250$ MeV) (full lines) together with fit according to eq. (6) (dashed line).

Since in the transport theoretical calculation described below the mean field is evaluated assuming a thermal momentum distribution, it is useful to parameterize the above potential in a simple form

$$U(p) = V_0 \left(1 - \left(\frac{p}{a_1} \right)^2 \right) \exp \left(- \left(\frac{p}{a_1} \right)^2 \right) \left(\frac{T}{T_0} \right)^3, \quad (6)$$

where V_0 is the value of the potential at zero momentum while a_1 corresponds to the momentum where the potential changes sign. The temperature dependence in this parameterization reflects the fact that the optical potential is essentially proportional to the pion density. In fig. 1 we compare this parameterization for three

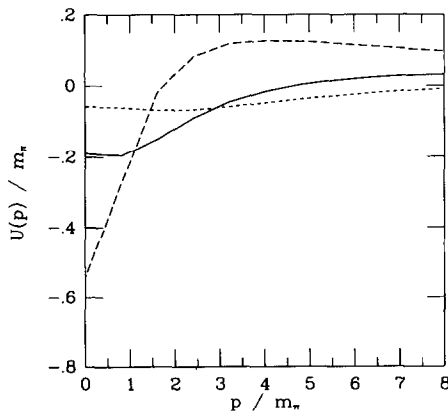


Fig. 2. Mean-field potential (3) for different resonance masses. Full line: $M_{\text{res}} = 500$ MeV; long-dashed line: $M_{\text{res}} = 770$ MeV; short-dashed line: $M_{\text{res}} = 1100$ MeV.

TABLE 1
Different sets of parameters of effective potential (6)

Parameter set	V_0 [MeV]	a_1 [MeV]	T_0 [MeV]
Set 1	40	650	200
Set 2	100	650	200
Set 3	100	1300	200

different temperatures ($T = 150, 200, 250$ MeV) with the potential based on $\pi + \pi \rightarrow \rho$. The parameters for the fit are $V_0 = -0.2m_\pi$, $a_1 = 4m_\pi$ and $T_0 = 200$ MeV. The agreement is reasonable over the whole range of temperatures displayed. For the sake of simplicity in the following we will only use the parameterization (6) for the mean-field potential. Also the meaning of the parameters becomes more transparent with this choice.

In the following we will use several parameter sets for the potential (6) which we display in table 1. Set 1 corresponds to the potential obtained by Shuryak for the pure pion gas. Since our purpose is a general understanding of mean-field effects, we have also calculated with other parameter sets shown in the table. Both these are much stronger than the realistic potential, in order to see the effects of a strong potential.

3. Static potential

Before we turn to the expansion of the fireball it is instructive to first study the simpler case of particles leaving a static potential well. Because the energy E of the particle is conserved while traversing the potential the *energy* spectra dN/dE inside and outside are identical,

$$\frac{dN}{dE_{\text{inside}}} = \frac{dN}{dE_{\text{outside}}}. \quad (7)$$

Let us assume that the particles inside the potential are distributed according to a boost-invariant fire tube,

$$\frac{dN}{dy d^2p_\perp} = \int_{-\infty}^{\infty} d\eta m_\perp \cosh(\eta) \exp(-\beta m_\perp \cosh(\eta)) = 2m_\perp K_1(\beta m_\perp), \quad (8)$$

where for simplicity we have assumed that the particles are distributed according to Boltzmann statistics ($\beta = 1/T$). Thus, in a given rapidity bin, the problem reduces to a two-dimensional one, with $E = m_\perp = \sqrt{m^2 + p_\perp^2}$.

Because of relation (7), the potential only affects the momentum spectra; for cylindrical symmetry we have

$$\frac{dN}{d^2p} = \frac{1}{p} \frac{dN}{dp} = \frac{1}{p} \frac{dE}{dp} \frac{dN}{dE}. \quad (9)$$

Outside the potential we have

$$\frac{dE}{dp} = \frac{p}{E}, \quad (10)$$

thus we find

$$\frac{dN}{d^2p_{\perp, \text{outside}}} = \frac{1}{E} \frac{dN}{dE_{\text{outside}}} = \frac{1}{E} \frac{dN}{dE_{\text{inside}}} = \frac{1}{E} \left(\frac{p}{dE/dp} \frac{dN}{d^2p_{\perp, \text{inside}}} \right). \quad (11)$$

Finally using eq. (8) for the momentum distribution inside the potential we find

$$\frac{dN}{dE_{\text{outside}}} = \frac{p_{\text{in}}}{(dE/dp)_{\text{in}}} \frac{dN}{d^2p_{\perp, \text{inside}}}, \quad (12)$$

where p_{in} is determined from

$$\sqrt{p_{\text{in}}^2 + m^2} + U(p_{\text{in}}) = E, \quad (13)$$

and

$$\frac{dE}{dp_{\text{in}}} = \frac{p_{\text{in}}}{\sqrt{p_{\text{in}}^2 + m^2}} + \frac{dU}{dp_{\text{in}}}. \quad (14)$$

In fig. 3 we have plotted the resulting spectra for the different parameterizations (table 1) of the potential together with the free one. For the spectra inside the potential we have used a temperature of $T = 1/\beta = 135$ MeV, which fits the transverse-momentum spectra from proton-proton collisions (see below). All spectra are normalized such that they have the same number of particles *inside* the potential. Therefore, in fig. 3 the integral over the free spectrum is larger than the integral over the ones with a potential, because in the latter case all particles with $E < m$ are bound inside the well.

We find that, as a result of the potential, the spectra become somewhat steeper at small momenta. The effect, however, is much too small in order to explain the data. Even if we increase the depth of the potential well from $V_0 = -40$ MeV (short-dashed line) to $V_0 = -100$ MeV (long-dashed line), the slope of the spectrum is not changed much (aside from an overall shift downwards due to the fact that the deeper potential keeps more particles bound inside). The difference between parameter set 2 (short-dashed line) and 3 (long-dashed line) is even smaller.

This behaviour actually can be understood by looking at fig. 4, where we have plotted the resulting dispersion relations $E(p)$ for the three parameterizations. Only particles with energy above the rest mass $E(p) \geq m$ contribute to the outside spectrum. There the slopes of the dispersion relation dE/dp do not differ very much from the free one. As a consequence the depth of the potential does not affect the spectra very much. Also possible minima of the dispersion relation^{14,15)} do not affect the outside spectrum, because they also occur at energies smaller than the rest mass.

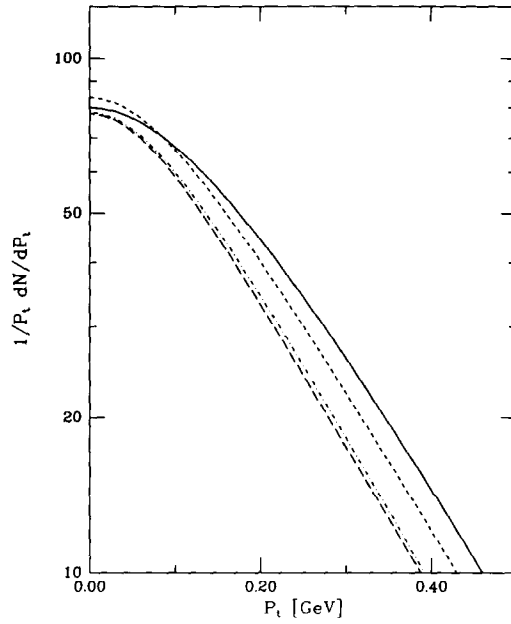


Fig. 3. Spectra for static potential based on eq. (6) and parameterizations of table 1. Free: full line; set 1: short-dashed line; set 2: long-dashed line; set 3: dashed-dotted line.

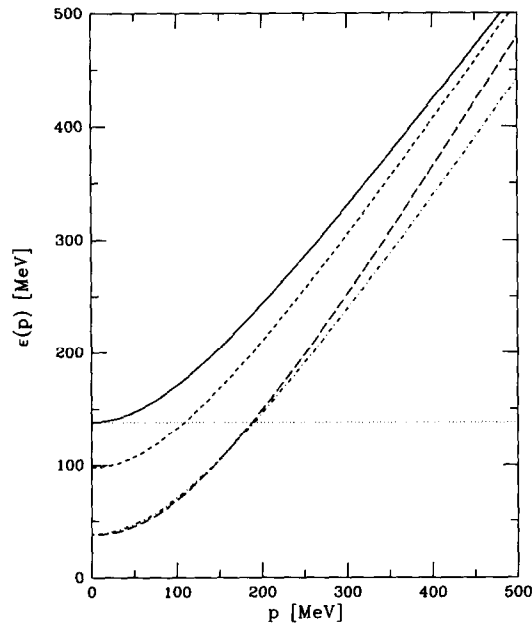


Fig. 4. Dispersion relation for parameterizations of table 1. Labels as in fig. 3.

In conclusion, we find, that in the (not very realistic) case of a static potential well, only small effects of the potential on the particle spectrum can be observed, too small to account for the observed enhancement. The main reason is that the particles which may escape the potential feel only a rather weak potential with very moderate momentum dependence.

Of course as already pointed out, in reality the potential is not static because its source, the fireball, expands. Consequently the potential decreases as time continues so that eventually all particles may escape. The observed enhancement could, therefore, still originate from those particles, which would be trapped inside the static potential.

This would be the case if these particles leave the system as soon as energy conservation allows the particles to do so, i.e. once their energy becomes larger than their rest mass. However, such a scenario, depends very much on the time scales involved in the problem. It would require that the expansion of the fireball is very slow compared to the velocity of the low-momentum particles. Since the fireball is made out of pions as well this seems to be very unlikely.

On the other hand, if the fireball expands with a velocity faster than that of the low-momentum particles these particles will essentially remain inside the potential well (fireball) until the potential has vanished. In this case only the fast particles would feel an effect of the potential as they have to climb a potential well of finite depth. At the soft part of the spectrum, however, we would not expect any effects of the potential*.

The question of the time scales involved can best be answered in a model calculation. In the next section we, therefore, will study the full expansion of the fireball and the effect of the collective potentials on the spectra in a transport model. This approach should provide a reasonable simulation of the expansion and the time scales involved.

4. Expansion of the fireball

In order to study the dynamic effects of the collective potential we have extended a cascade model¹⁹⁾ to include the propagation of the particles in a mean field. Following the standard procedure¹⁶⁾ the particle coordinates and momenta are propagated according to Newton's equations of motion,

$$\begin{aligned}\frac{dr}{dt} &= \frac{p}{E} + \frac{dU}{dp}, \\ \frac{dp}{dt} &= -\frac{dU}{dr},\end{aligned}\tag{15}$$

where the mean-field potential U is given by eq. (6). The temperature which is needed in order to determine the mean-field potential is calculated from the density

* This later scenario has been studied in great detail in case of anti-protons in ref. ¹⁸⁾.

of pions assuming local thermal equilibrium. The pion density is extracted from the actual particle distribution. At every time step the radial dependence of the density is fitted with a function of the form

$$\rho(r_{\perp}) = A(1 + b^2 r_{\perp}^2) \exp(-b^2 r_{\perp}^2). \quad (16)$$

The parameters A and b are determined by the root-mean-square radius of the distribution and the total number of pions. In the longitudinal direction, on the other hand, for a given time step the density is assumed to be constant between z_{\min} and z_{\max} , where z_{\min} and z_{\max} denote maximum distance in positive and negative direction where particles have materialized. Comparing with the actual density distributions these assumption are very well justified for the conditions we are dealing with and which we will discuss below*. The error introduced is certainly not larger than the one one would have when using a spatial grid and having to deal with large density fluctuations²⁰⁾.

The cascade includes π -, η -, ρ - and ω -mesons but no baryons**. All mesons decay according to their empirical lifetime and decay channels. For the pion-pion scattering measured phase shifts are used while for all other elastic processes a constant cross section of 20 mb is assumed. The only inelastic process taken into account is $\pi + \pi \leftrightarrow \rho$. In addition the ω is allowed to decay into three pions. While this may not take into account all the details of the meson-meson scattering it certainly provides enough accuracy in order to lead to a reasonably realistic expansion scheme. Finally the model does not respect the Bose nature of the mesons, i.e. we do not have any Bose enhancement factors incorporated in the collision integral as done e.g. in refs. ^{12,13)}. Here, we are rather interested in dynamical effects due to collective potentials.

We shall specifically be concerned with the central 200 GeV/A $^{16}\text{O} + \text{Au}$ data of the NA35 collaboration¹⁾. This experiment measures negative particles and does not identify the pions explicitly. Usually one assumes a $\sim 10\%$ admixture of kaons, electrons and anti-protons. For simplicity, however, we neglect this fact and assume all negatives to be pions.

The number of initial particles is determined such that the measured rapidity distribution of pions is reproduced. We further impose Bjorken-like²²⁾ initial conditions: rapidity y_{boost} and longitudinal coordinate of a local thermal distribution are uniquely related by

$$z = \tau_0 \sinh(y_{\text{boost}}), \quad (17)$$

where τ_0 may be considered a formation or hadronization time. In other words particles materialize, i.e. participate in the expansion, only after their proper time is larger than the hadronization time τ_0 . Because of eq. (17) this implies that in the

* This is essentially a result of the Bjorken initial conditions we have imposed here (see below).

** At CERN energies the ratio of protons to π^- is about 1:6 for S+S collisions²¹⁾.

c.m. system particles which are created at a large longitudinal distance from the center will materialize at a later time. The initial radial distribution is assumed to follow the density profile of the oxygen projectile.

We will study different ways of populating the initial momentum space. One possibility is to assume local thermal and chemical equilibrium so that with given density the initial local temperature and the admixture of higher resonances is fixed. In this approach the free parameters of the model are the initial number of particles and τ_0 . Those can be determined by fitting the measured rapidity distribution and slope of the high-energy part of the pion spectrum. The lowest hadronization time which still leads to acceptable transverse-momentum spectra and rapidity distributions is $\tau \approx 8 \text{ fm}/c$. The number of initial particles is 380.

The other procedure assumes that the momenta of all initial particles are distributed according to a slope parameter for pp collisions of $T_0 = 135 \text{ MeV}$. Their relative multiplicities are distributed proportional to their degeneracy, i.e. $\pi:\eta:\rho:\omega = 3:1:9:3$. These statistical weights are very close to what one obtains in string-fragmentation models²³). In order to reproduce the measured rapidity distribution we start with 280 particles in the initial state.

The resulting rapidity distributions of pions for both initialization schemes obtained after $30 \text{ fm}/c$ are shown together with the experimental data of the NA35 collaboration¹) in fig. 5. These results have been obtained with the pure cascade without potentials. However, since in our model the density and hence the potentials are assumed to be constant along the longitudinal direction, the rapidity distribution is not affected by the potentials.

Let us now turn to the effect of the collective potentials. In figs. 6 and 7 we show the resulting transverse-momentum spectra using the different initial conditions described above. Fig. 6 corresponds to the local-equilibrium initialization while fig. 7 represents the string-fragmentation picture. Different expansion schemes are studied:

- (1) The particles do not interact but may decay (full line).
- (2) Particles may collide but do not feel any mean-field forces (short-dashed line).
- (3) Particles collide and are subject to the mean-field forces (long-dashed line and dashed-dotted line).

We display only the results for parameter sets 1 and 2 of table 1, because the results obtained with parameter set 3 are indistinguishable from the one obtained with parameter set 2.

First let us point out the difference between the expansion with and without collisions (full and dashed histogram) in the low- p_t spectra exhibited in fig. 7. Since the initial conditions do not correspond to local chemical equilibrium the collisions are still very effective and provide a significant enhancement at low transverse momentum. This result should also be seen in connection with the findings of ref. ¹³), where the effect of Bose statistics in the collision integral is studied. Considering the fact that our results have been obtained without Bose statistics, we conclude

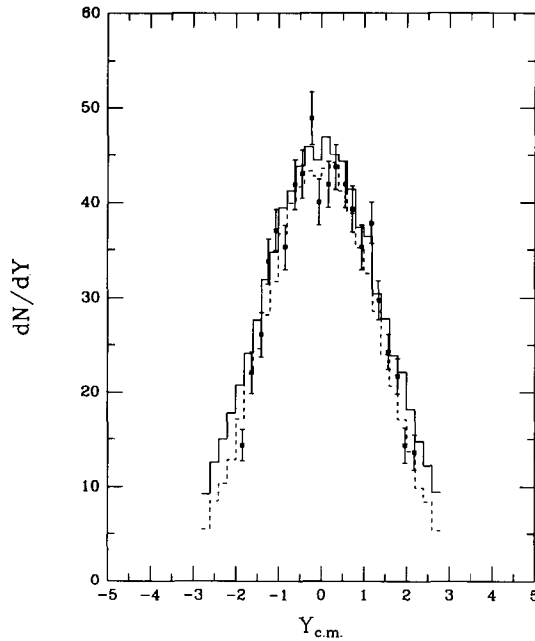


Fig. 5. Rapidity distribution for the two initialization schemes described in the text. Full histogram: local thermal equilibrium; dashed histogram: pp distribution.

that a great part of the effect demonstrated in ref. ¹³⁾ can be explained without any Bose statistics. This supports the finding pointed out in the aforementioned reference, namely that the effect of Bose statistics in the final state and in the matrix element cancel each other to a large extent. In case of the local-equilibrium initialization (fig. 6) the effect of the collision is very small. One of the reasons simply is that we had to choose a very high hadronization time τ_0 , in order to reproduce the high-momentum part of the spectra. As a consequence the initial density of particles is comparatively low and thus the collisions are not so effective.

For either initial conditions the mean field leads only to very little enhancement at low p_t . Also the stronger mean-field parameterization does not provide more enhancement. Thus, the expansion does not provide the additional enhancement compared to the static limit. Quite to the contrary, the effects found in the static limit are reduced because the time-averaged depth of the potential well is smaller as a result of the expansion. And certainly, the enhancement due to the pions being trapped in the static potential, as discussed in the last section, does not show up. Obviously, the velocity of the fireball expansion is larger than the velocity of the low-momentum pions. As already mentioned this is actually what one should expect, because the relevant part of the fireball, which provides the potential, consists of pions. The average velocity of these pions, however, is much larger than the one of the soft pions. Therefore, as long as there are no exotic transport phenomena, we

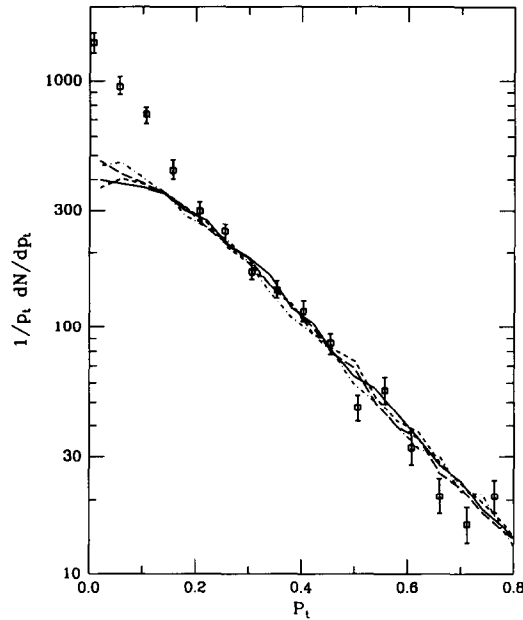


Fig. 6. Transverse-momentum spectra for local-equilibrium initialization: Full line: no mean field, no collision; short-dashed line: no mean field, collisions; long-dashed line: mean-field (set 1) collisions; dashed-dotted line: mean-field (set 2) collisions.

expect the expansion velocity of the source of the potential to be higher than the velocity of the low-momentum pions*. As a consequence the momentum of the pion remains essentially unchanged during the expansion. Thus, as long as we consider only one fluid (in our case the pions) an attractive potential does not lead to considerable enhancement in the spectrum at low momenta, because the velocity of the particles and the expansion velocity of the potential well are intimately related. This would be different if one had two different fluids, one which provides the potential, and one, the spectrum of which we want to study. In this case one could very well imagine that the expansion proceeds essentially adiabatically so that all bound pions would leave the potential with minimal momentum. Such a possibility actually exist at BEVALAC/SIS energy (~ 1 GeV) heavy-ion collisions. There the pions feel a collective potential through the coupling to the delta-hole channel²⁴), which looks very similar to the one we have studied here. In this case, however, the potential is provided by the nucleon fireball. Since the mass of the nucleon is considerably larger than the one of the pions, it could very well be that the expansion velocity of the nucleon fireball is smaller than the velocity of the low-momentum

* It is actually very difficult to imagine that an attractive two-particle interaction should slow down slow particles even more. Let us consider a pair of pions, fast and slow, which interact by an attractive potential. Clearly the effect of the potential is, to slow down the fast pion and to accelerate the slow one, such that both at the end have the same velocity.

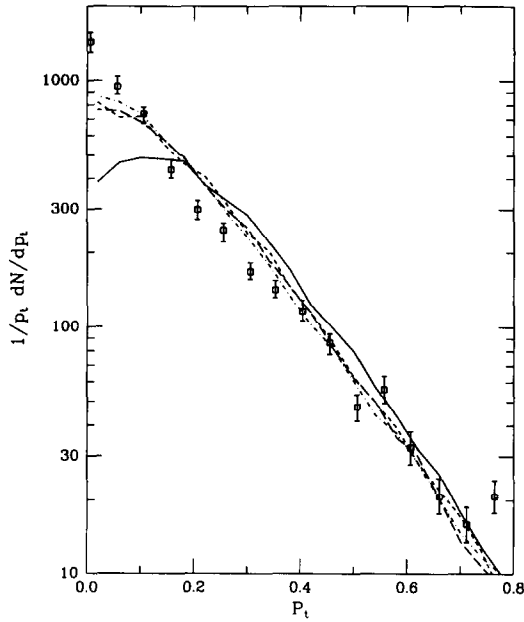


Fig. 7. Same as fig. 6, but with pp distribution in the initial state.

pions. Therefore, at BEVALAC/SIS energies the pion spectra could reveal information about the long sought in-medium pion dispersion relation*.

5. Conclusions

In conclusion, we could show that the observed enhancement at low transverse momentum in the spectra of pion cannot be accounted for by collective mean fields. We have pointed out, that this result depends very much on the time scales involved in the expansion. As long as the potential is provided by the same kind of particles as the ones the spectra are being studied of, we do not expect an enhancement, because in general the expansion velocity is faster than the velocity of the soft particles. This would be different at BEVALAC/SIS energies where the source for the pion collective potential would be given by the nucleons, which, as a result of their larger mass, would most likely expand slower than the soft pions.

In the case where the particles initially have been distributed according to the momentum spectrum of proton-proton collisions, we found that the inclusion of particle collision leads to considerable enhancement at low transverse momentum in comparison to the free decay, which would more or less correspond to the simple

* Indeed pionic spectra at these energies exhibit an enhancement at low energies similar to those discussed. This possibility will be addressed in a separate publication²⁵).

folding of the pp data. This effect, however does not fully account for the enhancement observed in the data.

Taking into account the results of ref. ¹³⁾ the soft-pion puzzle seems still unresolved. The data require either a very short hadronization time ($\tau_0 \approx 1 \text{ fm}/c$) or the absence of mesons heavier than the pion. Both assumptions certainly require a nontrivial scenario in order to be acceptable.

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