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PION MULTIPLICITY IN CENTRAL NUCLEUS–NUCLEUS COLLISIONS

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Abstract: Recently, the discrepancy between standard intranuclear cascade model predictions and experimental data for the pion multiplicity in central nucleus–nucleus collisions has been interpreted as a manifestation of compression effects. We explore here alternative effects, involving pion reabsorption by two nucleons, off-mass shell effects and renormalization of the pion sources inside a dense nuclear medium.

1. Introduction

The pion production rate in high-energy collisions between nuclei has recently attracted much interest and presents a puzzle. Various assumptions and models have been used to attempt an interpretation of the inclusive cross sections as well as the mean pion multiplicity from "central-trigger" measurements. In particular, Stock *et al.* [1]) have suggested that this latter quantity can bear information about the nuclear matter equation of state; their argumentation is based on some properties of the intranuclear cascade (INC) model [2,3]). To clarify the discussion, we briefly review the situation regarding pion multiplicity.

Current models of nucleus–nucleus collisions appear to overestimate the pion multiplicity. This is particularly true for the thermal model [4,5]) and for the phase-space model [6]), unless an unrealistically large freeze-out density is used. It is also true for the INC model, but the case is more complex here because specific dynamical assumptions are involved for pion production. Stimulated by the work of ref. [1]), our aim is to reexamine the INC model and to find out whether alternative mechanisms based on simple and plausible modifications of the model can shed light on the pion rate puzzle.

For collisions up to 2 GeV/A, it is generally assumed that pion production proceeds exclusively through the formation of Δ-isobars. This seems indeed reasonable in view of the predominance of the isobar in the collision of two nucleons in free space [7]). Since the emergence of pions is governed by the isobars, some uncertainty is introduced for nuclear pion production because of the poor knowledge of Δ-behaviour in a hot, dense nuclear medium. Assuming only elementary processes, the Δ's in the course of their propagation through nuclear matter may transform back into nucleons ($N\Delta \rightarrow NN$) and may also disappear by decay ($\Delta \rightarrow \pi N$); the released pions may form Δ's again ($\pi N \rightarrow \Delta$). While the cross sections and the

lifetime are known in free space, the medium may affect them significantly. For instance, in cold nuclei the Pauli blocking effect can increase the decay lifetime by about a factor two [8]).

Faced with uncertainties on the behaviour in the nuclear environment, we have previously studied [2,3]) two extreme cases:

(i) The Δ's have a long lifetime (i.e. longer than the nucleus–nucleus collision time; their decay is thus frozen during the whole collision process and the pions are created only at the end).

(ii) The Δ-lifetime keeps its natural value, and the pion yield therefore results from the interplay of Δ-production (in NN and πN collisions), recombination and decay.

Case (i) gives a spectrum in poor agreement with the data, displaying an unwanted hole in the so-called mid-rapidity region, but the predicted pion yield compares reasonably with experiment, although it is larger. In case (ii) the right spectral shape is given, but the model predicts too many pions, especially near the effective threshold ($\sim$400 MeV/A). Let us recall that these results have been obtained with a model accounting for pion reabsorption through the two-step ($\pi N \to \Delta$, $N\Delta \to NN$) process.

The discrepancy between the INC predictions and experiment was interpeted in ref. [1]) as a signal of the role of the compressional energy, during the stage when the pions are created. Close examination of the cascade evolution shows that the final pion population is obtained rather early in the collision, when maximum compression exists in the system. The observed discrepancy for the pion yield can be translated into the compression energy, i.e. the amount of energy necessary to reduce the kinetic (or thermal) energy, which is assumed to directly determine the observed pion multiplicity. In other words, the authors of ref. [1]) postulate that the INC model "makes the cascades too hot", and the heat which has to be removed to bring agreement with the data is transformed into compressional energy; the latter is then associated with the density reached in the cascade calculation. Of course, such a procedure may appear questionable since it relies on the properties of the INC model to obtain a quantitative energy-density relation which is absent from the model itself. At any rate, the specificity of the assumption and the plausible form of the result of ref. [1]) (a quadratic dependence) have the merit of making prominent the question of the proper pion-production mechanism inside the INC model.

That is the problem we want to reexamine in this paper. We will try to answer the following question: In such a context can we modify the ingredients and hopefully improve the pion production mechanism and the predictions for the pion yield?

More precisely, we study here three different lines of approach: (i) we introduce into our INC code the – so-far neglected – absorption of pions involving two nucleons at the same time; (ii) we investigate the influence of dealing with nucleons off their mass or energy shell; (iii) finally, we examine a (crude) way of introducing medium corrections for the reactions producing pions.

The paper is organized as follows. In sect. 2, we briefly describe the main features of our INC code. In sect. 3, we describe the implementation of the pion absorption by two nucleons and present our results for the Ar+KCl case. Sect. 4 is devoted to a study of the effect of the nucleons being off their mass shell. This situation is dictated by the binding energy of the nuclei, and has been neglected so far. We shortly discuss the problems tied to the implementation of such an effect in a simple classical model such as the INC. Numerical results are presented and shown to go in the desired direction, thus providing a possible clue for the understanding of pion production. In sect. 5, we tackle in a brief way the question of the intensity of the pion source inside the nuclear medium, and we tentatively propose a way to mock it up in the frame of the cascade model. Finally, sect. 6 presents the discussion of our results and the conclusion.

2. The intranuclear cascade model

Our INC model is described in refs. [2,3]). We briefly recall here the main points for the sake of completeness.

The nucleon positions are distributed randomly inside two touching spheres: projectile and target. In each nucleus, the nucleons are given momenta selected at random according to the Fermi gas law. The projectile is then Lorentz-boosted against the target. Nucleons travel at constant speeds until, for a pair of nucleons a and b, the minimum relative distance d_{ab} is reached. If d_{ab} is larger than $(\sigma_{tot}/\pi)^{1/2}$ the motion of the nucleons does not change, but if $d_{ab} \leq (\sigma_{tot}/\pi)^{1/2}$ the momentum vectors of the nucleons a and b are changed in agreement with energy and momentum conservation laws and according to the distribution law given by the experimental differential cross section. When reaction (or decay) channels are open, the nature of the particles is also changed. The collision is assumed to take a zero duration and straight-line motion is resumed until another pair reaches its minimum relative distance, and so on, until the collision rate becomes negligible. The procedure is repeated a sufficient number of times to reach satisfactory statistics. The observables result from ensemble averaging. Isospin degeneracy is assumed and potential wells are neglected so far.

We have included the following reactions:

$$\mathrm{NN} \to \mathrm{NN}\,, \qquad \mathrm{N}\Delta \rightleftarrows \mathrm{NN}\,, \qquad \mathrm{N}\Delta \to \mathrm{N}\Delta\,, \qquad \Delta\Delta \to \Delta\Delta\,, \qquad \pi\mathrm{N} \rightleftarrows \Delta\,. \tag{2.1}$$

The NN→NΔ process is assumed to account for the inelastic nucleon–nucleon scattering. This seems to be a reasonable assumption up to a nucleon kinetic energy ~2 GeV in the lab system. In this scheme pion production as well as pion annihilation are two-step processes. The cross section for the πN→Δ reaction is taken as the isospin average pion–nucleon cross section, a simple Breit–Wigner form with central energy $E_0 = 1232$ MeV and width $\Gamma = 115$ MeV inside the $(E_0-\Gamma, E_0+\Gamma)$ interval. The cross section for NΔ→NN is obtained from NN→NΔ by detailed balance.

This model gives a satisfactory agreement with experiment for the inclusive and doubly inclusive proton cross sections [3]). It gives reasonable pion spectra but, as stated in the introduction, predicts too large π^- multiplicities.

To sum up, the INC model is a purely classical model (except for the probabilitistic construction of the final configuration in elementary collisions) which builds up the nuclear process as the succession of *binary* collisions (and Δ-decays) between particles *on their mass shell*, proceeding as in *free space*. We are now going to examine how the model can be extended beyond these three (italicized) conditions, in relation to the pion multiplicities.

3. Two-nucleon absorption

One possible way out of the problem is trying to increase the pion absorption. Keeping the dynamical input as given by the free-space cross section leaves few possibilities. One of these is the often-invoked absorption of a pion by a pair of nucleons, a mechanism which can be represented as

$$\pi NN \to NN \,, \tag{3.1}$$

which is to be contrasted with the two-step process

$$\pi N \to \Delta \,, \qquad \Delta N \to NN \tag{3.2}$$

cited above. The importance of process (3.1) is very hard to determine from true pion absorption inside nuclei [9,10]). Even the observation of a back-to-back emission of two nucleons is not helpful since it is consistent with both (3.1) and (3.2). Only the energy dependence of the correlation yield may be decisive. As far as we know, the only nucleus for which these experimental data exist is the deuteron. They are given in fig. 1, which is adapted from ref. [9]). The huge bump around $T_\pi \approx 140$ MeV is due to the formation of a Δ-resonance. A similar bump dominates the total cross section. The shift observed in the absorption cross section is not well understood yet. Nevertheless, the low-energy part of the cross section seems to have an origin different from the resonant capture mechanism and suggests the presence of a direct $\pi NN \to NN$ process. The resonant pion annihilation by two nucleons is automatically included in the INC model as described in the previous section (except, of course, for possible quantum interferences). To avoid double counting, we thus have to include only the non-resonant direct contribution, which amounts to 5 mb at low T_π. We have adopted such a value for values of T_π up to 200 MeV. It should not be excluded after all that this mechanism extends so high in energy and that the shift of the resonant peak is due to interference of the amplitudes for the two mechanisms.

Technically, the $\pi NN \to NN$ process is implemented in the INC code as follows. Pion–nucleon pairs are checked when they reach their minimum relative distance $d_{\min}$. If the c.m. pair energy $\sqrt{s}$ is larger than 1120 MeV (the energy of the

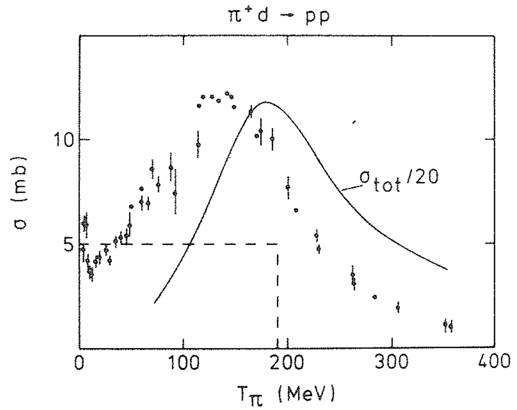

Fig. 1. Experimental $\pi^+ d \to pp$ cross section (dots), adapted from ref. [9]). The full curve represents the total $\pi^+ + d$ cross section divided by 20. The dotted line is the cross section, adopted in our calculation (see text), for the direct ($\pi NN \to NN$) absorption by two nucleons. The quantity T_π is the kinetic pion energy in the lab system.

Δ-resonance minus the width), the quantity $\pi d^2_{\min}$ is compared to the $\pi N \to \Delta$ cross section, as previously. Below this value, the closest nucleon neighbour is searched, and the distance d separating the pion from the c.m. of the two nucleons is calculated. The quantity πd^2 is compared to the non-resonant $\pi NN \to NN$ cross section as defined above in order to decide whether such a process occurs. The results of such a calculation are shown in fig. 2. The effect of the two-nucleon absorption is very small except at 400 MeV beam energy, where it amounts to 20%. It is more important at low energy because the average pion kinetic energy is lower. In conclusion, the effect of the $\pi NN \to NN$ process is quite small and its magnitude is not large enough to remove the discrepancy. We recall however that this discussion is based solely on the deuteron data.

4. Off-shell effects

We here approach the possibility of having particles off their mass shell. In the standard treatment, the interaction energy is completely neglected, assuming everything is dominated by the kinetic energy of the particles. In view of the good success of the model around 1 GeV/A, this seems legitimate in this energy range. Around 400 MeV/A, corrections due to the interactions are likely to be more important. The potential energy of the nucleons within a nucleus becomes significant as compared to the incident energy per nucleon in the c.m. system. Interaction effects due to interacting pairs may be very intricate. As a first step, we adopt a quasi-particle picture, with particles off their energy shell. Even in this simple picture, we have very few pieces of knowledge which might provide a clue for an appropriate procedure.

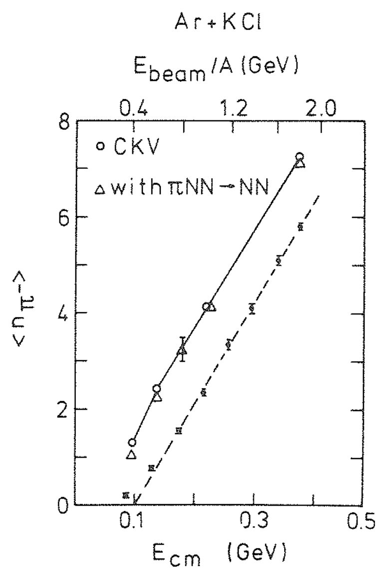

Fig. 2. Excitation function for the pion multiplicity in central collisions of Ar+KCl. The experimental data (full dots) correspond to the "central trigger" [see ref. [18])]. The open dots refer to the standard INC calculations of ref. [2]). The triangles give the results when the $\pi NN \rightarrow NN$ process is included into the calculation. We have indicated by an error bar the typical uncertainty of the calculations (also valid for figs. 3 and 4).

In a very simple-minded picture, based on nonrelativistic concepts, one may consider that the nucleons are initially moving in potential wells, providing a binding field in the individual nuclei. If we want to keep on with relativistic kinematics a new difficulty arises, linked with the several ways of introducing a potential in the relativistic context. We may have a scalar field with the following energy-momentum relation

$$E^2 = p^2 + (m + V_s)^2 , \qquad (4.1)$$

or a vector field, with

$$(E - V_v)^2 = p^2 + m^2 . \qquad (4.2)$$

Much attention has been recently paid to the relativistic description of the nuclear mean field [for a review, see ref. [11])]. It seems that, within such a picture, both an attractive scalar and a repulsive vector field should be introduced, and each of them is a few hundreds of MeV strong. In the following, we have decided to consider a field of one type only at a time, either vector or scalar, to see whether the nature of the field itself plays a role.

An even more delicate problem concerns the dynamic evolution of the field as the nucleus–nucleus collision proceeds. Very likely, in a violent explosive collision, the average field is ultimately destroyed, but there is no indication at which rate the baryon field is decaying. We have therefore considered two extreme cases:

(a) The field survives the collision process, or, in other words, the particles remain off the energy shell all the time.

(b) Once the particles have made a collision, they do not experience the binding field anymore.

The latter case pictures the highest rate of destruction of the field. We note, however, that this assumption is not unreasonable, since collisions scatter nucleons in the momentum space. Even if they stay within the volume of the incoming nuclei, they are kicked far apart from the other nucleons in phase space [12,13]). It is plausible to consider that the fields are localized in the regions of phase space occupied by the incoming nucleons. This is roughly consistent with the empirically known energy dependence of the real part of the nucleon–nucleus optical potential [14]).

To simplify the calculation, we have considered a uniform potential, either scalar or vector. Its magnitude can be adjusted to reproduce the correct available energy in the initial state. A value V_s or V_v around -40 MeV is required. The time dependence is chosen as in (a) and (b) described above. Since these correspond to extreme conditions, our results should be interpreted as upper and lower limits.

Before presenting the results, we have to elaborate on the off-energy shell cross sections. For the integrated $(1+2\rightarrow 3+4)$ cross section, one may write

$$\sigma \propto |T(s)|^2 R = \pi |T(s)|^2 \frac{p_{\text{c.m.}}}{\sqrt{s}}, \tag{4.3}$$

where T is the average (over final states) transition matrix and where R is the available two-body phase space in the exit channel. The quantity $p_{\text{c.m.}}$ is the outgoing c.m. momentum. The cross section (4.3) is directly given by an experimental measurement, if the particles are taken on their mass shell. We now want to relate the cross section for off mass shell particles to an on the mass shell cross section. Let us consider particles 1 and 2 having a c.m. momentum p. For the on-mass shell case, the invariant c.m. energy is

$$\sqrt{s_0} = (p^2+m_1^2)^{1/2} + (p^2+m_2^2)^{1/2}, \tag{4.4}$$

and for the off-mass shell case (for a scalar potential)

$$\sqrt{s} = [p^2+(m_1+V_s)^2]^{1/2} + [p^2+(m_2+V_s)^2]^{1/2}. \tag{4.5}$$

For simplicity, we consider the same potential for all the particles. In first order of V_s,

$$\sqrt{s} = \sqrt{s_0} + 2\frac{m_1+m_2}{\sqrt{s_0}} V_s. \tag{4.6}$$

Now, let us turn to $p_{\text{c.m.}}$. In the free case, we have

$$p^0_{\text{c.m.}} = \frac{[(s_0-(m_3+m_4)^2)(s_0-(m_3-m_4)^2)]^{1/2}}{2\sqrt{s_0}}, \tag{4.7}$$

whereas, for 3 and 4 lying in a potential V_s, we have

$$p_{\text{c.m.}}=[(s-(m_3+m_4+2V_s)^2)(s-(m_3-m_4)^2)]^{1/2}/2\sqrt{s}\,. \tag{4.8}$$

Taking account of (4.6), and using $(m_3-m_4)^2 \ll s$, which is reasonable in our case, we get

$$\frac{p_{\text{c.m.}}}{\sqrt{s}}\approx\frac{1}{\sqrt{s_0}}\left[(s_0-(m_3+m_4)^2)\left(1-4\frac{(m_1+m_2)}{s_0}V_s\right)+4(m_1+m_2-m_3-m_4)V_s\right]^{1/2}. \tag{4.9}$$

For an elastic collision, this relation becomes

$$\frac{p_{\text{c.m.}}}{\sqrt{s}}\approx\frac{p^0_{\text{c.m.}}}{\sqrt{s_0}}\left(1-2\frac{(m_1+m_2)V_s}{s_0}\right). \tag{4.10}$$

With $V_s\approx-40$ MeV, and for nucleon masses, the correction is less than 4%. For an inelastic collision, the last term in the bracket of eq. (4.9) introduces in addition a small shift to the effective threshold. For the production of a Δ-particle, this shift is of the order of 5 MeV. We have neglected these small effects, and used an off-mass shell cross section $\tilde{\sigma}(\sqrt{s})$ of the form

$$\tilde{\sigma}(\sqrt{s})\propto|T(s)|^2\frac{p^0_{\text{c.m.}}}{\sqrt{s_0}}, \tag{4.11}$$

or, since T is usually a smooth function of s,

$$\tilde{\sigma}(\sqrt{s})\approx\sigma(\sqrt{s_0})\approx\sigma\left(\sqrt{s}-2\frac{m_1+m_2}{\sqrt{s}}V_s\right), \tag{4.12}$$

where σ is the on-mass shell cross section. These considerations show that in case (a) mentioned above, the cross sections are not practically changed by the introduction of a potential, provided the incident momentum vectors are kept fixed.

As for case (b), we evidently have

$$\tilde{\sigma}(\sqrt{s})=\sigma(\sqrt{s}), \tag{4.13}$$

since the final state momentum is the same as in the free case. If a vector potential V_v is used, the considerations above are qualitatively valid. It is easy to check that in the case (a), we have

$$\tilde{\sigma}(\sqrt{s})\approx\sigma(\sqrt{s}-2V_v) \tag{4.14}$$

instead of (4.12). In the case (b), the same equation as (4.13) is obtained.

From the practical point of view, we have considered attractive potentials V_s and V_v of 40 MeV depth. We have used the same potential for nucleons and Δ's. For the latter, a smaller potential would probably be better, but we kept the same value for simplicity. The pions are produced on their mass shell. We will come back to

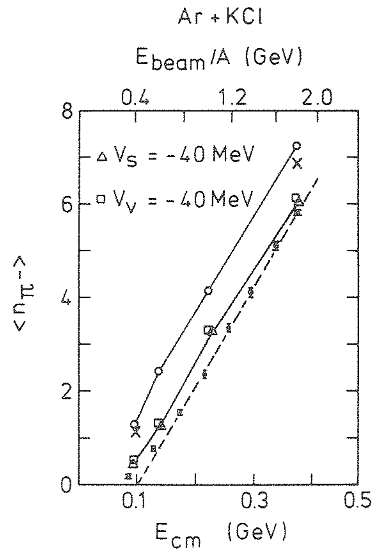

Fig. 3. Same as fig. 2. The open triangles show the result when the nucleons are put off their mass shell by a constant attractive scalar potential of 40 MeV depth and when the time evolution is governed by prescription (b) (see text). The open squares correspond to a constant attractive vector potential of same depth with the same prescription. The crosses indicate schematically the results when the scalar potential is used with prescription (a) (no destruction of the potential, see text).

this point in the discussion. The same exponential dependence upon t (the four-momentum transfer) as in the free case is adopted for the elastic NN differential cross section.

The results are presented in fig. 3. The triangles and the squares indicate the results for case (b) (with rapid destruction of the field). They are almost in agreement with experimental data. As expected, the results with case (a) (no destruction of the field) are very close to the no potential case (open dots). They are only sketched in the figure for the sake of clarity. The difference amounts to 5 to 10%. These results by no means solve the puzzle, but, according to us, indicate that the maximum effect is of the order of the original discrepancy.

Let us finally notice that the momentum distribution of the nucleons is not sensitively changed by the introduction of the off-shell effects. The particles are slightly less dispersed in the perpendicular direction: the off-shell effects (in case (b)) decrease somewhat the number of binary collisions.

5. Quenching of the pion sources

Nucleons (possibly excited) can be considered as sources for the pion field [15]. Inside a nucleus, the strength of the sources can be modified because of the presence of the surrounding matter, which polarizes the source. This problem has been treated repeatedly during the past ten years [15,16]) in the static limit. Retaining here the presentation of ref. [15]), the free space classical field equation for the pion field ϕ,

$$(-\nabla \cdot \nabla + m_\pi^2)\phi = \frac{g_r}{2m_N} \nabla \cdot \sigma \,, \tag{5.1}$$

where g_r is the pion-nucleon coupling constant, and where σ is the nucleon spin, is to be modified into

$$-\nabla \cdot [1+\alpha_e(x)]\nabla\phi + m_\pi^2\phi = -\frac{g_r}{2m_N} \nabla \cdot \eta(x)\sigma(x) \,. \tag{5.2}$$

The medium has two effects. First, it modifies the pion propagator by a kind of dielectric constant $\varepsilon = 1+\alpha_e(x)$, which arises from the axial polarizability of the nucleons by the field itself. This question has been studied extensively, even for real pions, in connection with pion condensation [17]). The second modification is the renormalization of the strong-interaction coupling constant g_r by the factor η,

$$\eta = 1 + \tfrac{1}{3}\alpha_e \,. \tag{5.3}$$

The quantity α_e is given by

$$\alpha_e = \frac{\alpha}{1-\frac{1}{3}\alpha} \,, \tag{5.4}$$

where α is the polarizability per unit volume of the nuclear medium. It is proportional to the baryon density. The renormalization of the sources arises from the short-range correlations, which repel the nucleons from each other and, in a sense, prevent the pion field propagating from its source. The factor $\frac{1}{3}$ comes from a simplified treatment of the short-range correlations between nucleons, similar to the Lorentz–Lorenz counterpart in electromagnetism. Recent developments indicate that the factor $\frac{1}{3}$ is reduced when a proper treatment is used [17]). A complete field-theoretic treatment of the pion production, including these effects, in the case of relativistic nucleus–nucleus collisions is not feasible. The question we address here is thus the following: How can we introduce such an effect in a cascade calculation? If we limit ourselves to indicative results, a possible way is to multiply the pion creation (more exactly, $\mathrm{NN} \to \mathrm{N}\Delta$) cross section by the function η defined above. This procedure is rather crude, but could give a rough idea of the order of magnitude of such an effect on the pion yield. We applied the same correction to the $\mathrm{N}\Delta \to \mathrm{NN}$ cross section.

In practice, we re-ran the calculation, looking at the baryon density around the colliding NN pairs. To avoid fluctuations, we took the average density $\bar{\rho}$ in a sphere of 2 fm radius positioned at the c.m. of the colliding pairs. We just modified the $\mathrm{NN} \to \mathrm{N}\Delta$ cross section by the factor

$$\eta = 1 - \tfrac{1}{2}A\bar{\rho} \,. \tag{5.5}$$

With the values quoted in ref. [15]), $A \approx 4.4\ \mathrm{fm}^3$, which makes a reduction of about 20% in normal nuclear matter. The results are given in fig. 4. The pion yield is

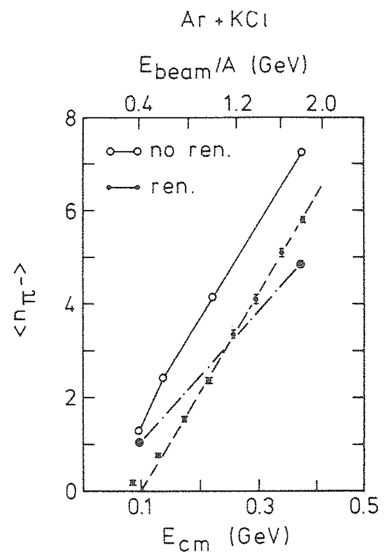

Fig. 4. Same as fig. 2. The dot-and-dashed line shows the result when the NN⇄NΔ cross sections are multiplied by a density-dependent function η (see text), which corresponds to a renormalization of the pion source strength due to their immersion in a dense medium.

reduced as expected, but too much at high energy and not enough around 400 MeV/A. This, of course, reflects the larger densities reached at high energy. Nevertheless, our results show that the correction due to the quenching of the sources could be of the order of the discrepancy between the standard calculation and the experiment.

6. Discussion

We took again the problem of the discrepancy between standard INC predictions and experimental data concerning the pion yield in central nucleus–nucleus collisions in the 400 MeV/A–2 GeV/A range. Our purpose was to find out whether the INC predictions could be improved without calling for compressional energy effects as in ref. [1]).

We first looked at the possible absorption of pions by two nucleons at the same time, without exciting one of these to a delta resonance. This process reduces the pion yield, but by a small (insufficient) amount, if one uses experimental information about pion capture by deuterons. Of course, one may wonder whether the cross section would increase when the correlated pair is in a more compact state as it is presumably in the fireball created during the nucleus–nucleus collisions; but no experimental data exist, which could bring information on this aspect.

Afterwards, we investigated the effect of putting nucleons off their mass shell as is mandatory, owing to the fact that nuclei are bound systems. We underlined that there are many ways to put particles off their mass shell and we have studied two of them, just for illustration, since the proper method for handling this situation is not known. The main uncertainty, however, seems to be tied to the evolution of the potential in the course of the nuclear collision. If it is assumed that the potential survives throughout the collision, the pion production rate is not much affected, as could have been anticipated. On the other hand, the field may be rapidly destroyed; we have treated such a situation via the specific assumption that a particle ceases to feel any potential as soon as it has scattered for the first time. In the latter case, the pion yield is dramatically decreased and is close to experimental values. However, this result must be considered as an upper limit for the effect, because of the ambiguities due to the time dependence as well as the energy dependence (that we have neglected) of the potential ascribed to mock up off-mass shell effects. Furthermore, we have neglected off-shell effects for pions, which might be important, because there is no simple way to include them.

Finally, we investigated through a simplified procedure the effect of the possible quenching of pion field sources when immersed in a nuclear medium. Our input was based on the extensive study of the zero-frequency classical pion field in nuclei. It turns out that this effect lowers the pion yield, has the right order of magnitude, but a wrong energy dependence. However, many points have to be clarified. First, it is not well known how the pion sources are renormalized in the finite frequency (real pion) regime. Second, and more importantly, it is not known to what extent a renormalization of the sources can be simply translated in a renormalization of the delta production cross section, as we have assumed.

In conclusion, we do not pretend to have shown the absence of compressional energy effects, but we think we have shown that, if the latter exist, they are intermingled with other effects, namely off-mass shell effects and renormalization of the pion sources, whose calculation is however subject to many uncertainties. We would finally like to note that the pion production mechanism might be so complex (see sect. 5) that its full treatment goes beyond the scope of a simple model like INC and requires a field-theoretic approach.

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