Can pions created in high-energy heavy-ion collisions produce a Centauro-type effect?

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Abstract

We study a Centauro-type phenomenon in high-energy heavy-ion collisions by assuming that pions are produced semiclassically both directly and in pairs through the isovector channel.

The leading-particle effect and the factorization property of the scattering amplitude in the impact-parameter space are used to define the classical pion field. By analyzing the joint probability function $P_{II_3}(n_0,n_-)$ for producing $n_0$ neutral and $n_-$ negative pions from a definite isospin state $II_3$ of the incoming leading-particle system we show that only direct production of pions without isovector pairs favors Centauro-type behavior. The presence of isovector pairs seems to destroy the effect. Our conclusion is supported through the calculation of two pion correlation parameters, $f_2^{0-}$ and $f_2^{00}$, and the average number of neutral pions $\langle n_0 \rangle_{n_-}$ as a function of negative pions $n_-$ produced.

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In recent years several cosmic-ray experiments [1] have reported evidence for the existence of Centauro events characterized by an anomalously large number of charged pions in comparison with the number of neutral pions, indicating that there should exist a very strong long-range correlation between two types of the pions. Such long-range correlations are possible if pions are produced coherently and constrained by the global conservation of isospin [2–7].

Although the actual dynamical mechanism of the production of a classical pion field in the course of a high-energy collisions is not known, there exist a number of interesting theoretical speculations [8–13] that localized regions of misaligned vacuum might occur in ultrahigh-energy hadronic and heavy-ion collisions. These regions become coherent sources of a classical pion field. In early models [2,3], however, the coherent production of pions is taken for granted and is considered to be a dominant mechanism.

These models also predict strong negative correlations between the number of neutral and charged pions. In fact, the exact conservation of isospin in a pion uncorrelated jet model is known [4,5] to give the same pattern of neutral/charged fluctuations as observed in Centauro events. This strong negative neutral-charged correlation is believed to be a general property of the direct pion emission in which the cluster formation (or the short-range correlation between pions) is not taken into account [6,7,14].

A similar conclusion may be drawn when a proper multipion symmetrization combined with the isospin singlet channel is considered [15].

In this paper we consider the leading-particle effect as a possible source of a classical pion field [16]. Pions are assumed to be produced from a definite isospin state of the incoming leading-particle system both directly and in isovector pairs.
Coherently emitted isovector clusters decay subsequently into pions outside the region of interaction. We discuss the limiting behavior of the probability distribution of neutral pions, the neutral - to - charged ratio and the corresponding two-pion correlation function \( f^{ab}_2 \), as well as the variation of the average number of neutral pions \( \langle n_0 \rangle \) as a function of the number of negative pions \( n_- \) produced.

We show that the Centauro effect is strongly suppressed if pions are produced in isovector pairs. At high energies most of the pions are produced in the central region. To isolate the central production, we adopt high-energy longitudinally dominated kinematics, with two leading particles retaining a large fraction of their incident momenta. The energy available for the hadron production is

\[
E_{had} = \frac{1}{2} \sqrt{s} - E_{leading}. \tag{1}
\]

It is related to the inelasticity \( K \) of the collisions by

\[
E_{had} = \frac{1}{2} \sqrt{s} K. \tag{2}
\]

Note that at fixed total c.m. energy \( \sqrt{s} \), the hadronic energy \( E_{had} \) varies from event to event.

With a set of independent variables \( s, \{ \vec{q}_iT, y_i \} \equiv q_i, \ i = 1, 2, \ldots n \), the n-pion contribution to the s-channel unitarity becomes an integral over the relative impact parameter \( b \) of the two incident leading particles:

\[
F_n(s) = \frac{1}{4s} \int d^2b \prod_{i=1}^n dq_i \ | T_n(s, \vec{b}; 1 \ldots n) |^2, \tag{3}
\]

where \( dq = d^2q_Tdy/2(2\pi)^3 \). The normalization is such that

\[
F_n(s) = s\sigma_n(s), \quad \sigma_{inel}(s) = \sum_{n=1}^{\infty} \sigma_n(s). \tag{4}
\]
At high energies the matter distributions of the two-leading-particle system are Lorentz-contracted disks in the c.m. system. We assume that, in the early stages of the collision, a highly excited localized system occurs that relaxes through the coherent emission of pions [17]. According to [18], this type of coherent emission of pions should saturate close to the threshold since, once it starts, the resulting pion emission, in the absence of a "resonance cavity", prevents further buildup of pion fields. The basic assumption of the independent pion-emission model, neglecting the isospin for a moment, is the factorization of the scattering amplitude $T_n(s, \vec{b}; 1 \ldots n)$ in the $b$ space:

$$T_n(s, \vec{b}; 1 \ldots n) = 2s f(s, \vec{b}) \frac{n^{-1}}{\sqrt{n!}} \prod_{i=1}^{n} J(s, \vec{b}; q_i),$$

where, owing to unitarity,

$$| f(s, \vec{b}) |^2 = e^{-\bar{n}(s, \vec{b})},$$

and

$$\bar{n}(s, \vec{b}) = \int dq | J(s, \vec{b}; q) |^2$$

denotes the average number of emitted pions at a given impact parameter $b$. The function $| J(s, \vec{b}; q) |^2$, after the integration over $b$, controls the shape of the single-particle inclusive distribution. A suitable choice of this function also guarantees that the energy and the momentum are conserved on the average during the collision.

The above particular factorization of $T_n$ was considered earlier by R. Avin et al. [19] in connection with the construction of a unitary model of multiparticle production.

The inclusion of isospin in this model is straightforward [3]. We observe that the factorization of $T_n$ is a consequence of the pion field satisfying the equation of motion

$$\left(\Box + \mu^2\right) \vec{\pi}(s, \vec{b}; x) = \vec{j}(s, \vec{b}; x),$$

where $\mu$ is the mass of the pion.
where $\vec{j}$ is a classical source related to $\vec{J}(s, \vec{b}; q)$ via the Fourier transform

$$\vec{J}(s, \vec{b}; q) = \int d^4xe^{i\vec{q}\cdot\vec{x}}\vec{j}(s, \vec{b}; x).$$

(9)

The standard solution of Eq.(8) is given in terms of in- and out-fields that are connected by the unitary $S$ matrix $\hat{S}(\vec{b}, s)$ as follows:

$$\vec{\pi}_{\text{out}} = \hat{S}^\dagger \vec{\pi}_{\text{in}} \hat{S} = \vec{\pi}_{\text{in}} + \vec{\pi}_{\text{classical}},$$

(10)

where

$$\vec{\pi}_{\text{classical}} = \int d^4x'\Delta(x - x'; \mu)\vec{j}(s, \vec{b}; x').$$

(11)

The $S$ matrix following from such a classical source is still an operator in the space of pions. Inclusion of isospin requires $\hat{S}(s, \vec{b})$ to be also a matrix in the isospace of the leading particles.

The coherent production of isovector clusters of pions is described by the following $S$ matrix:

$$\hat{S}(s, \vec{b}) = \int d^2\vec{e} |\vec{e}\rangle D(\sum_c \vec{J}_c; s, \vec{b})\langle \vec{e}|,$$

(12)

where $|\vec{e}\rangle$ represents the isospin-state vector of the two-leading-particle system. The quantity $D(\sum_c \vec{J}_c; s, \vec{b})$ is the unitary coherent-state displacement operator defined as

$$D(\sum_c \vec{J}_c; s, \vec{b}) = \exp[\sum_c \int dq \vec{J}_c(s, \vec{b}; q)a_c^\dagger(q) - H.c.],$$

(13)

where $a_c^\dagger(q)$ is the creation operator of a cluster $c$ and the summation $\sum_c$ is over all clusters. The clusters decay independently into $c = 1, 2, \ldots$ pions outside the region of strong interactions.

Clusters decaying into two or more pions simulate a short-range correlation between pions and behave as $\vec{\rho}, \vec{\rho}', \ldots$, etc. in isospace. They need not be well-defined resonances. The more pions in a cluster, the larger the correlation effect.
expected. If the conservation of isospin is a global property of the colliding system, restricted only by the relation
\[ \vec{I} = \vec{I} + \vec{I}_{\pi}, \]
(14)
where \( \vec{I}_{\pi} \) denotes the isospin of the emitted pion cloud, then \( \vec{J}_c(s, \vec{b}; q) \) is of the form
\[ \vec{J}_c(s, \vec{b}; q) = J_c(s, \vec{b}; q) \vec{e}, \]
(15)
where \( \vec{e} \) is a fixed unit vector in isospace independent of \( q \). The global conservation of isospin thus introduces the long-range correlation between the emitted pions.

If the isospin of the system of two incoming (outgoing) leading particles is \( II_3 \) and \( I' I'_3 \), respectively, then the initial-state vector of the pion field is \( \hat{S}(s, \vec{b}) | II_3 \rangle \), where \( | II_3 \rangle \) is a vacuum state with no pions but with two leading particles in the isostate characterized by \( II_3 \). Then the \( n \)-pion production amplitude is
\[ iT_n(s, \vec{b}; q_1 \ldots q_n) = 2s \langle I' I'_3; q_1 \ldots q_n | \hat{S}(s, \vec{b}) | II_3 \rangle. \]
(16)
The unnormalized probability distribution of producing \( n_+ \pi^+ \), \( n_- \pi^- \), and \( n_0 \pi^0 \) pions is defined as
\[ W(n_+ n_0, n'_0, II_3; II'_3) = \int d^2b dq_1 dq_2 \ldots dq_n \langle II'_3; n_+ n_0 | \hat{S}(s, \vec{b}) | II_3 \rangle |^2, \]
(17)
where \( n = n_+ + n_- + n_0. \)

Assuming further that all \( (I', I'_3) \) are produced with equal probability, we can sum over all possible isospin states of the outgoing leading particles to obtain
\[ P_{II_3}(n_+ n_0) = \frac{\sum_{II'_3} W(n_+ n_0, II'_3; II_3)}{\sum_{n_+ n_0} \sum_{II'_3} W(n_+ n_0, II'_3; II_3)}. \]
(18)
This is our basic relation for calculating various pion-multiplicity distributions, pion
multiplicities, and pion correlations between definite charge combinations.

In order to obtain some more detailed results for multiplicity distributions and correlations, one should have an explicit form for the source function $J_c(s, \vec{b}; q)$, $c = 1, 2, \ldots$.

The results on the isospin structure are most easily analyzed in the so-called grey-disk model in which

$$\pi_c(s, \vec{b}) = \pi_c(s) \theta(b_0(s) - b),$$

where $\pi_c(s)$ denotes the mean number of clusters of the type $c$, $b_0(s)$ is related to the total inelastic cross section, and $\theta$ is a step function.

Let us assume that pions are produced both directly and through isovector clusters of the $\rho$ type. In this case, the distribution of neutral pions becomes

$$P_{II_3}(n_0) = (I + 1/2)(I - I_3)! \int_{-1}^{1} dx |P_{I_3}^{I_3}(x)|^2 \frac{[A(x)]^{n_0}}{n_0!} e^{-A(x)},$$

where $A(x) = \pi_\pi x^2 + (1 - x^2)\pi_\rho$. Here $\pi_\pi$ denotes the average number of directly produced pions, and $\pi_\rho$ denotes the average number of $\rho$-type clusters which decay into two short-range correlated pions. The function $P_{I_3}^{I_3}(x)$ denotes the associate Legendre polynomial.

Note that the total number of emitted pions is

$$\langle n \rangle = \pi_\pi + 2\pi_\rho.$$

The correlation $f_2$ between the pions,

$$f_2 = \langle n(n - 1) \rangle - \langle n \rangle^2 = 2\pi_\rho,$$

4In general, the probability $W(n, n, n_0, I'I_3'; II_3)$ depends on $(I'I_3')$ dynamically. The final-leading-particle production mechanism usually tends to favor the $(I'_3', I_3') \approx (I, I_3)$ case, for example, if the final-leading particles are nucleons or isobars. However, if the leading particles are colliding nuclei, it is reasonable to assume almost equal probability for various $(I'_3', I_3')$ owing to the large number of possible leading isobars in the final state.
is positive indicating deviation from the Poisson distribution \((f_2 = 0)\) owing to the emission of \(\rho\)-type clusters.

In Fig. 1 we show the behavior of \(P_{II_3}(n_0)\) for \(\langle n \rangle = 50\) and different combinations of \((\pi, \rho)\) when (a) \(I = I_3 = 0\), and (b) \(I = I_3 = 1\). We see that Centauro-type behavior is obtained only for pions that are produced directly, i.e., for \(\pi \neq 0\) and \(\rho = 0\).

If \(R = \frac{n_0}{\langle n \rangle}\) denotes the fraction of neutral pions, then it is easy to see that in the limit \(\langle n \rangle \to \infty\) with \(R\) fixed, the probability distribution \(\langle n \rangle P_{II_3}(n_0)\) scales to the limiting behavior:

\[
\langle n \rangle P_{II_3}(n_0) \to P_{II_3}(R) = (I + \frac{1}{2}) \frac{(I - I_3)!}{(I + I_3)!} \left| P_{II_3}^{I_3}(x_R) \right|^2 \left[ \left( \frac{R}{I_3} \right) \left( \frac{1 - \gamma_{\rho}(R - \gamma_{\rho})}{I_3} \right) \right]^{-\frac{1}{2}},
\]

where \(\gamma_{\rho} = \frac{\pi_{\rho}}{\langle n \rangle} = \frac{\pi_{\rho}}{\langle n \rangle + 2\pi_{\rho}}\) and

\[
x_R^2 = \frac{R - \gamma_{\rho}}{1 - 3\gamma_{\rho}}, \quad 0 \leq \gamma_{\rho} < \frac{1}{3}.
\]

This limiting probability distribution is different from the usual Gaussian random distribution for which one expects peaking at \(R = \frac{1}{3}\) as \(\langle n \rangle \to \infty\).

Instead, we find peaking at \(R = \gamma_{\rho}\) although

\[
\langle n_0 \rangle_{II} = \langle n \rangle - \langle n_{ch} \rangle_{II} = \frac{1}{2I + 3} \frac{[\pi_{\pi} + 2(I + 1)\pi_{\rho}]}{\langle n_{ch} \rangle_{II}}
\]

and \(\langle n_0 \rangle = \frac{1}{2} \langle n_{ch} \rangle = \frac{1}{3} \langle n \rangle\), for either \(I = 0\) or \(\pi = \rho\).

If \(\gamma_{\rho} = \frac{1}{3}\), that is if \(\pi_{\pi} = \pi_{\rho}\), the probability \(P_{II_3}(n_0)\) becomes a pure Poisson distribution and is independent of \(II_3\).
In Fig. 2 we show the average number of neutral pions \( \langle n_0 \rangle_{\pi_-} \), as a function of the number of negative pions \( n_- \) produced for different pairs of \( (\pi_\pi, \pi_\rho) \) and \( I = I_3 = 1 \). We note that \( \langle n_0 \rangle_{\pi_-} \) decreases with \( n_- \) only in the region where \( \pi_\pi \gg \pi_\rho \) and \( \langle n \rangle \gg n_- \), indicating again that the Centauro-type effect cannot be expected if a significant number of \( \rho \)-type clusters of pions is produced. Similarly we find that the neutral-pion two-particle correlation parameters \( f_a^{00} = \langle n_0 n_a \rangle - \langle n_0 \rangle \langle n_a \rangle - \langle n_0 \rangle \delta_{0a} \) for \( a = 0 \) and \( - \) are related as

\[
f_2^{00} + \frac{1}{2} f_2^{00} = \pi_\rho \frac{I + 1}{2I + 3}, \tag{26}
\]

where

\[
f_2^{00} = \frac{4(I + 1)}{(2I + 3)^2(2I + 5)}(\pi_\pi - \pi_\rho)^2, \tag{27}
\]

showing that \( f_2^{00} \) is negative if \( \pi_\rho \) is small. We should also mention earlier works on neutral-pion correlation effects [20].

In summary, the results of the present analysis show that the emission of isovector clusters of pions in the framework of a unitary eikonal model with the global conservation of isospin strongly suppresses the Centauro-type effect for pions. This might suggest that the Centauro-type effect, if any, will probably appear only in very limited regions of phase space where isovector clusters should be missing. How it is possible dynamically is not yet clear [10–13]. Similar discussion about emission of isoscalar clusters of pions is presented in [15].

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Figure captions:

Fig. 1. Multiplicity distributions $P_{II_3}(n_0)$ of neutral pions for $\langle n \rangle = 50$ when the total isospin of the incoming-leading-particle system is (a) $I = I_3 = 0$ and (b) $I = I_3 = 1$. The curves represent different combinations of $(n_\pi, n_\rho)$, the average number of directly produced pions, and the average number of $\rho$-type clusters, respectively.

Fig. 2. The average number of neutral pions as a function of the number of negative pions produced for $I = I_3 = 1$. The curves represent different combinations of $(\bar{n}_\pi, \bar{n}_\rho)$. 
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