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

M. Martinis,* V. Mikuta-Martinis,† and J. Črnugelj‡

Department of Physics, Theory Division, Rudjer Bošković Institute, P.O.B. 1016, 41001 Zagreb, Croatia

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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. We show that the Centauro-type effect is strongly suppressed if a large number of pions are produced in isovector pairs. Our conclusion is supported through the calculation of two-pion correlation parameters, f_2^{0-} and f_2^{00}, as well as f_{2,n_-}^0 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 the two types of 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 collision is not known, there exist a number of interesting theoretical speculations [8–13] that localized domains in which the chiral condensate is disoriented might occur in relativistic hadronic and heavy-ion collisions. These domains become coherent sources of a classical pion field. However, current studies of formation of these domains show that they are in fact too small (essentially pion sized) for the coherent production of pions to be seen experimentally [11]. In early models [2,3] 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 is combined with a model where all pions arise from emission of isoscalar pairs [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 in the central region from a definite isospin state of the incoming leading-particle system both directly and in isovector pairs.

These 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_2^{ab}, as well as the variation of f_{2,n_-}^0 and the average number of neutral pions ($\langle n_0\rangle_{n_-}$) as a function of the number of negative pions (n_-) produced.

We show that the Centauro-type effect is strongly suppressed if a large number of pions is 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

$$E_{\rm had} = \frac{1}{2}\sqrt{s} - E_{\rm leading} \ . \tag{1}$$

It is related to the inelasticity K of the collisions by

$$E_{\rm had} = \frac{1}{2}\sqrt{s}K \ . \tag{2}$$

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

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

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

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

$$F_n(s) = s\sigma_n(s) \ ,$$
$$\sigma_{\rm inel}(s) = \sum_{n=1}^{\infty} \sigma_n(s) \ . \tag{4}$$

*Electronic address: martinis @thphys.irb.hr
†Electronic address: mikutama@thphys.irb.hr
‡Electronic address: crnugelj@thphys.irb.hr

At sufficiently high colliding energies the matter distributions of the two-leading-particle system (nuclei) become Lorentz-contracted disks in the c.m. system that pass through each other and deposit their kinetic energy gradually behind them. We assume that, in the early stages of the collision, the colliding particles form a highly excited localized classical system that relaxes through the coherent emission of pions [17] and their clusters. 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 and clustering for a moment, is the factorization of the scattering amplitude $T_n(s,\vec{b};1\cdots n)$ in the b space:

$$T_n(s,\vec{b};1\cdots n)=2sf(s,\vec{b})\frac{i^{n-1}}{\sqrt{n!}}\prod_{i=1}^{n}J(s,\vec{b};q_i)\ , \tag{5}$$

where, owing to unitarity,

$$|f(s,\vec{b})|^2=e^{-\overline{n}(s,\vec{b})}\ , \tag{6}$$

and

$$\overline{n}(s,\vec{b})=\int dq|J(s,\vec{b};q)|^2 \tag{7}$$

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 Aviv *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 in the impact-parameter space of the leading-particle system

$$(\Box+\mu^2)\vec{\pi}(s,\vec{b};x)=\vec{j}(s,\vec{b};x)\ , \tag{8}$$

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

$$\vec{J}(s,\vec{b};q)=\int d^4x\,e^{iqx}\vec{j}(s,\vec{b}x)\ . \tag{9}$$

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:

$$\vec{\pi}_{\rm out}=\hat{S}^{\dagger}\vec{\pi}_{\rm in}\hat{S}=\vec{\pi}_{\rm in}+\vec{\pi}_{\rm classical}\ , \tag{10}$$

where

$$\vec{\pi}_{\rm classical}=\int d^4x'\Delta(x-x';\mu)\vec{j}(s,\vec{b};x')\ . \tag{11}$$

The S matrix following from such a classical course is 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:

$$\hat{S}(s,\vec{b})=\int d^2\vec{e}|\vec{e}\rangle D\left(\sum_c\vec{J}_c;s,\vec{b}\right)\langle\vec{e}|\ , \tag{12}$$

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

$$D\left(\sum_c\vec{J}_c;s,\vec{b}\right)=\exp\left[\sum_c\int dq\vec{J}_c(s,\vec{b};q)\vec{a}_c^{\dagger}(q)-\text{H.c.}\right]\ , \tag{13}$$

where $\vec{a}_c^{\dagger}(q)$ 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, which means that the final-state interaction between pions is neglected.

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

$$\vec{I}=\vec{I}'+\vec{I}_{\pi}\ , \tag{14}$$

where $\vec{I}_{\pi}$ denotes the isospin of the emitted pion cloud, then $\vec{J}_c(s,\vec{b};q)$ should be of the form

$$\vec{J}_c(s,\vec{b};q)=J_c(s,\vec{b};q)\vec{e}\ , \tag{15}$$

where $\vec{e}$ is a fixed unit vector in isopsace 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. The n-pion production amplitude is

$$iT_n(s,\vec{b};q_1\cdots q_n)=2s\langle I'I'_3;q_1\cdots q_n|\hat{S}(s,\vec{b})|II_3\rangle\ . \tag{16}$$

With $\hat{S}(a,\vec{b})$ of the form (12) we obtain

$$\langle I'I'_3;q_1\cdots q_n|\hat{S}(s,\vec{b})|II_3\rangle=\sum_{I_\pi=|I-I'|}^{I+I'}C^{(I_\pi)}(I'I'_3;II_3)\times\langle q_1\cdots q_n|I_\pi I_{\pi3};s,\vec{b}\rangle_{\rm coh} \tag{17}$$

with $I_{\pi 3} = I_3 - I'_3$, where

$$C^{(I_\pi)}(I'I'_3; II_3) = \int d^2e\, Y^*_{I'I'_3}(\vec{e}) Y_{I_\pi I_{\pi 3}}(\vec{e}) Y_{II_3}(\vec{e}) \quad (18)$$

and

$$|I_\pi, I_{\pi 3}; s, \vec{b}\rangle_{\rm coh} = \int d^2e\, Y_{I_\pi I_{\pi 3}}(\vec{e}) \prod_c D(J_c\vec{e}; s, \vec{b})|0\rangle \ . \quad (19)$$

The unnormalized probability distribution of producing $n_+\pi^+$, $n_-\pi^-$, and $n_0\pi^0$ pions is defined as

$$W(n_+n_-n_0, I'I'_3, II_3) = \int d^2b\, dq_1 dq_2 \cdots dq_n \times |\langle I'I'_3 n_+n_-n_0|\hat{S}(s,\vec{b})|II_3\rangle|^2 \ , \quad (20)$$

where $n = n_+ + n_- + n_0$.

If all the isospins (I', I'_3) of the outgoing-leading-particle system are produced with equal probability, then we can sum over all $(I'I'_3)$ using group theory alone to obtain

$$P_{II_3}(n_+n_-n_0) = \frac{\sum_{I'I'_3} W(n_+n_-n_0, I'I'_3; II_3)}{\sum_{n_+n_-n_0} \sum_{I'I'_3} W(n_+n_-n_0, I'I'_3; II_3)} \ . \quad (21)$$

This is now our basic relation for calculating various pion-multiplicity distributions, pion multiplicities, and pion correlations between definite charge combinations. In 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 $(I', I'_3) \approx (I, I_3)$ case, for example, if the final leading particles are nucleons or isobars. However, if the leading particles are colliding nuclei, an almost equal probability for various (I', I'_3) seems a reasonable approximation owing to the large number of possible leading isobars in the final state. In our model Eq. (12) the dynamical information which would restrict the possible values $(I'I'_3)$ is not present. Nevertheless, even if we assume that $(I', I'_3) \approx (I, I_3)$ the cloud of pions clusters could still have the isospin $I_\pi = 0, 2, ..., 2I$.

Recent studies of heavy-ion collisions at the partonic level [20] argue that the central region is mainly dominated by gluon jets. The valence quarks of the incoming nuclei together with the gluons which escape from the interaction region form the outgoing-leading-particle system. Since the gluon's isospin is zero, it is very likely that the total isospin of the produced pions in the central region is also zero. This picture is certainly true if the central region is free from valence quarks, the situation expected to appear at extremely high collision energies. In our model the case $I_\pi = 0$ follows if we set $I = I' = 0$ in Eq. (17) and $I = 0$ in Eq. (21).

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

The isospin structure of our model is most easily analyzed in the so-called gray-disk model in which

$$\overline{n}_c(s, \vec{b}) = \overline{n}_c(s)\theta(b_0(s) - b) \ , \quad (22)$$

where $\overline{n}_c(s)$ denotes the mean number of clusters of the type c, $b_0(s)$ is related to the total inelastic cross section, and θ is a step function. We assume that pions are produced both directly and through isovector clusters of the ρ type.

In order to facilitate the study of two-pion correlation functions we define the generating function $G_{II_3}(z, n_-)$ as

$$G_{II_3}(z, n_-) = \sum_{n_0, n_+} P_{II_3}(n_+, n_-, n_0) z^{n_0} \ , \quad (23)$$

from which we calculate

$$\langle n_0 \rangle_{n_-} = \frac{d}{dz} \ln G_{II_3}(1, n_-) \ , \quad (24)$$

$$f^0_{2,n_-} = \frac{d^2}{dz^2} \ln G_{II_3}(1, n_-) \ , \quad (25)$$

and

$$P_{II_3}(n_0) = \frac{1}{n_0!} \frac{d^{n_0}}{dz^{n_0}} \sum_{n_-} G_{II_3}(0, n_-) \ .$$

The form of the generating function $G_{II_3}(z, n_-)$ is the following:

$$G_{II_3}(z, n_-) = (I + \tfrac{1}{2}) \frac{(I - I_3)!}{(I + I_3)!} \int_{-1}^{1} dx |P_I^{I_3}(x)|^2 \times \frac{[A(z,x)]^{n_0}}{n_0!} e^{-B(z,x)} \ , \quad (26)$$

where

$$2A(z,x) = (1 - x^2)\overline{n}_\pi + z(1 - x^2)\overline{n}_\rho + 2x^2\overline{n}_\rho \quad (27)$$

and

$$2B(z,x) = \overline{n}_\pi(1 + x^2 - 2zx^2) + \overline{n}_\rho[2 - z(1 - x^2)] \ . \quad (28)$$

Here $\overline{n}_\pi$ denotes the average number of directly produced pions, and $\overline{n}_\rho$ denotes the average number of ρ-type clusters which decay into two short-range correlated pions. The function $P_I^{I_3}(x)$ denotes the associated Legendre polynomial. Note that $A(1, x) = B(1, x)$.

The total number of emitted pions is

$$\langle n \rangle = \overline{n}_\pi + 2\overline{n}_\rho \ . \quad (29)$$

The correlation f_2 between the pions,

$$f_2 = \langle n(n-1) \rangle - \langle n \rangle^2 = 2\overline{n}_\rho \ , \quad (30)$$

is positive, indicating deviation from the Poisson distribution ($f_2 = 0$) owing to the emission of ρ-type clusters.

In Fig. 1 we show the behavior of $P_{II_3}(n_0)$ for $\langle n \rangle =$

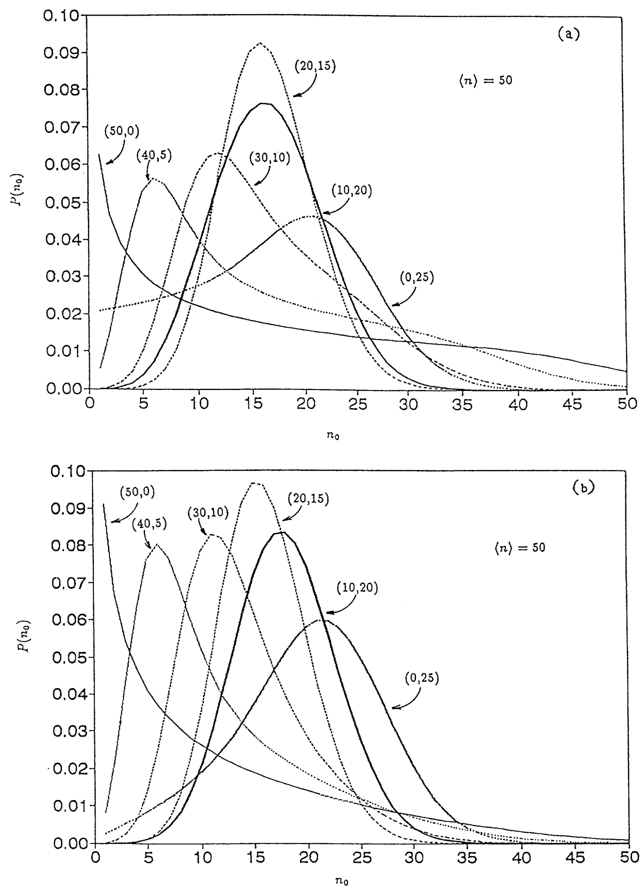

FIG. 1. Multiplicity distributions $P_{II_3}(n_0)$ of neutral pions for $\langle n\rangle = 50$ when the total isospin of the incoming-leading-particle system is (a) $I = I_3 = 0$ and (b) $I = I_3 = 1$. The curves represent different combinations of $(\overline{n}_\pi, \overline{n}_\rho)$, the average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

50 and different combinations of $(\overline{n}_\pi, \overline{n}_\rho)$ 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 $\overline{n}_\pi \neq 0$ and $\overline{n}_\rho = 0$.

In Fig. 2 we show the behavior of $\langle n_0\rangle_{n_-}$ for different combinations of $(\overline{n}_\pi, \overline{n}_\rho)$ when $I = I_3 = 1$. We note that $\langle n_0\rangle_{n_-}$ decreases with n_- only in the region where $\overline{n}_\pi \gg \overline{n}_\rho$ and $\langle n\rangle \gg n_-$, indicating again that the Centauro-type effect cannot be expected if a significant number of ρ-type clusters of pions is produced. Recent estimate of the ratio of ρ mesons to pions at accelerator energies is $\overline{n}_\rho = 0.10\overline{n}_\pi$.

In Fig. 3 we show the behavior of the dispersion $D(n_0)_{n_-}$ which is related to f^0_{2,n_-} as

$$f^0_{2,n_-} = [D(n_0)_{n_-}]^2 - \langle n_0\rangle_{n_-}$$

again for different pairs of $(\overline{n}_\pi, \overline{n}_\rho)$ and $I = I_3 = 1$. This quantity is believed to be sensitive to the pairing properties of the pions. The σ- and the ρ-type clusters behave differently.

If $R = n_0/\langle n\rangle$ denotes the fraction of neutral pions,

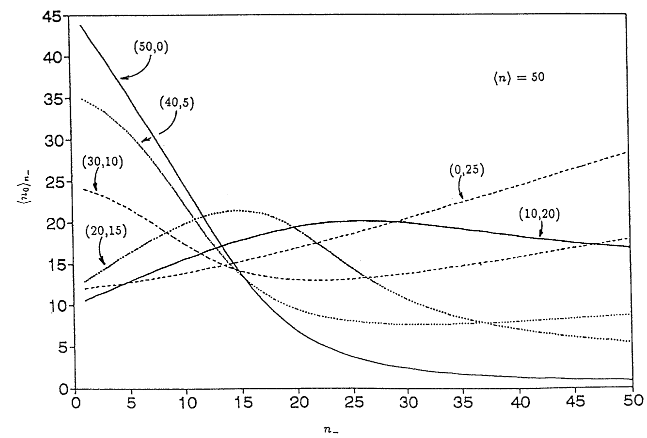

FIG. 2. The average number of neutral pions as a function of the number of negative pions produced for $I = I_3 = 1$. The curves represent different combinations of $(\overline{n}_\pi, \overline{n}_\rho)$.

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

$$\begin{aligned}\langle n\rangle P_{II_3}(n_0) &\to P_{II_3}(R)\\ &= (I + \tfrac{1}{2})\frac{(I - I_3)!}{(I + I_3)!}|P_I^{I_3}(x_R)|^2\\ &\quad\times[(1 - 3\gamma_\rho)(R - \gamma_\rho)]^{-1/2}\ , \end{aligned} \tag{31}$$

where $\gamma_\rho = \overline{n}_\rho/\langle n\rangle = \overline{n}_\rho/(\overline{n}_\pi + 2\overline{n}_\rho)$ and

$$x_R^2 = \frac{R - \gamma_\rho}{1 - 3\gamma_\rho}\ , \quad 0 \leq \gamma_\rho < \tfrac{1}{3}\ . \tag{32}$$

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

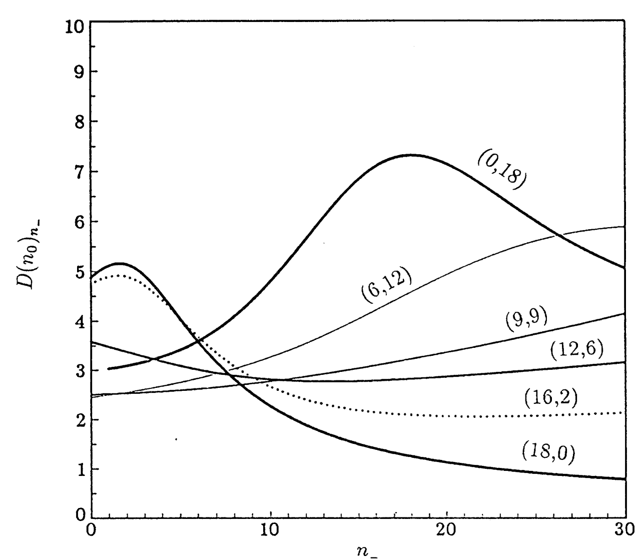

FIG. 3. The correlation function for two neutral pions as a function of number of negative pions for $I = I_3 = 1$. The curves represent different combinations of $(\overline{n}_\pi, \overline{n}_\rho)$.

Instead, we find peaking at $R = \gamma_\rho$ although

$$\langle n_0\rangle_{II} = \langle n\rangle - \langle n_{\rm ch}\rangle_{II} = \frac{1}{2I+3}[\overline{n}_\pi + 2(I+1)\overline{n}_\rho] \qquad (33)$$

and

$$\langle n_0\rangle = \tfrac{1}{2}\langle n_{\rm ch}\rangle = \tfrac{1}{3}\langle n\rangle ,$$

for either $I = 0$ or $\overline{n}_\pi = \overline{n}_\rho$. If $\gamma_\rho = \frac{1}{3}$, that is, if $\overline{n}_\pi = \overline{n}_\rho$, the probability $P_{II_3}(n_0)$ becomes a pure Poisson distribution and is independent of II_3.

Similarly, we find that the neutral-pion two-particle correlation parameters $f_2^{0a} = \langle n_0 n_a\rangle - \langle n_0\rangle\langle n_a\rangle - \langle n_0\rangle\delta_{0a}$ for $a = 0$ and $-$ are related in our model:

$$f_{2,II}^{0-} + \tfrac{1}{2}f_{2,II}^{00} = \overline{n}_\rho \frac{I+1}{2I+3} , \qquad (34)$$

where

$$f_{2,II}^{00} = \frac{4(I+1)}{(2I+3)^2(2I+5)}(\overline{n}_\pi - \overline{n}_\rho)^2 , \qquad (35)$$

showing that $f_{2,II}^{0-}$ is negative if $\overline{n}_\rho$ is small. We should also mention earlier works on neutral-pion correlation effects [21], due to independent cluster production, restricted by charge conservation.

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 [1,11–13]. A similar discussion about emission of isoscalar clusters of pions is presented in [15].

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