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DISORIENTED CHIRAL CONDENSATE AND CHARGE-NEUTRAL PARTICLE FLUCTUATIONS IN HEAVY-ION COLLISIONS

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The possibility of large charge and isospin fluctuations of the disoriented chiral condensate type (DCC-type) in high-energy heavy-ion collisions is studied within the framework of an unitary eikonal model. The factorization property of the scattering amplitude in the impact-parameter space of the leading two-nucleon system is used to study semiclassical production of pions in the central region. A classical source function of a pion field in the impact-parameter space was related to the semiclassical solutions of the equation of motion of the nonlinear σ-model coupled to quark degrees of freedom. A multi-pion exchange potential between two quarks is derived. It is shown that in the limit of soft chiral pion bremsstrahlung, the anomalously large fluctuations in the ratio of neutral to charged pions are obtained without involving DCC formation. We also show that the DCC-type fluctuations should be suppressed if a large number of pions was produced via ρ-type clusters.

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1. *Introduction*

Central ultrarelativistic collisions at RHIC and LHC with more then 3000 produced particles present remarkable opportunity to analyse event-by-event fluctuations of hadronic observables. Such single-event analysis with large statistics can reveal new physical phenomena usually hidden when averages over a large statistical sample of events are made [1, 2]. The number of particles produced in relativistic heavy-ion collisions can differ dramatically from collision to collision due to the variation of impact parameter (centrality dependence), energy deposition (leading particle effect) and other dynamical effects [2, 3]. The fluctuations can also be influenced by novel phenomena such as the disoriented chiral condensate (DCC) formation [4, 5].

The best probes of such novel dynamics are fluctuations of conserved quantities, because conservation laws limit the degree to which final-state scattering can dissipate. Even globally conserved quantities, such as energy, net charge, isospin, baryon number and strangeness can fluctuate when measured, e.g., in a limited phase-space region.

Several methods have been proposed [3, 6] to distinguish between statistical and dynamic fluctuations. In high-energy hadronic and heavy-ion collisions, the correlations of particle production are usually analyzed and described in terms of "short-range correlations" (SRC) and "long-range correlations" (LRC). The SRC are mainly due to clustering of outgoing particles. These clusters may be hadronic resonances which decay into a few particles or minijets [5, 7, 8]. The LRC dominate in high multiplicity events and are mainly due to global conservation laws.

The old puzzle in cosmic-ray observations is the existence of few exotic events characterized by an anomalously large number of charged pions in comparison with the number of neutral pions, the Centauros [9], indicating that there should exist a strong negative long-range correlation between two types of the pions. Such long-range correlations are possible if pions are produced semiclassicaly and constrained by global conservation of isospin [10 − 15].

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 numerous interesting recent theoretical attempts to explain Centauros either as different types of isospin fluctuations due to the formation of a DCC [16 − 20], or as multiparticle Bose-Einstein correlations (BEC) [21], or as the formation of a strange quark matter (SQM) [22]. Among the most interesting speculations is the idea of DCC that localized regions of misaligned chiral vacuum might occur during the ultrahigh-energy hadronic and heavy-ion collisions when the chiral symmetry is restored at high temperatures. These regions, if produced, would behave as a pion laser, relaxing to the ground state by coherent pion emission. It is generally accepted that the fluctuation of the ratio of neutral to charged pions of the Centauro type could be a sign of the DCC formation provided that a single large domain is formed, containing a large number of low p_{T} pions. Since the pions formed in the DCC are essentially classical, they form a quantum superposition of coherent states with different orientation in isospin space. If all pions in the domain are pointing in the same isospin direction and the condensate state is a pure isoscalar, then the formation of DCC leads to large event-by-event fluctuations in the ratio $f = n_0/n$ of the number of π_0's in the DCC divided by the total number of pions produced in an event. The probability distribution of f inside the DCC domain is [13, 16, 17, 19]

$$P_{\mathrm{DCC}}(f) = \frac{1}{2\sqrt{f}} \,. \tag{1}$$

There is a variety of proposed mechanisms other than DCC which also lead to the distribution (1) [23 − 26]. The distribution $P_{\mathrm{DCC}}(f)$ is different from the generic binomial-distribution expected in normal events which assumes equal probability for production of π_+, π_- and π_0 pions. The emission of charged and neutral pions is then uncorrelated and the distribution

$$P_{\rm B}(n_0, n) = \binom{n}{n_0} \left(\frac{1}{3}\right)^{n_0} \left(\frac{2}{3}\right)^{n-n_0} \tag{2}$$

in the limit $n \to \infty, n_0 \to \infty$ with f fixed, approaches a delta function at $f = 1/3$.

The possibility of observing the DCC-type fluctuations critically depends on the size and the energy content of the DCC domain. If the domain is of the pion size, the effect of DCC is too small to be observed experimentally. The early accelerator searches for Centauros and DCC at CERN [27 − 30] and at Fermilab [23, 31] were thus unsuccessful. With the RHIC facility at BNL now, there is a possibility to consider event-by-event fluctuations of other specific hadronic observables that might be more informative about the signal of DCC. There are various factors that may affect the observation of the DCC signal [32]. Experimental signals such as the isospin fluctuations, the strong relative enhancement of the number of low $p_{\rm T}$ pions, and the suppression of HBT correlations may provide a robust signal of DCC [33].

The space-time scenario of the formation and decay of DCC is usually studied within one of the simplified versions of the chiral effective Lagrangians, either the linear or nonlinear sigma model [20]. However, it should be emphasized that the use of σ-models, be they linear or nonlinear, is only a rough approximation to the true dynamics, because the couplings of pion and sigma fields to the constituent quarks may be large and their effect should not be ignored.

In this paper, we present results of an event-by-event analysis of charge-neutral pion fluctuations as a function of the ρ/π ratio in pp-collisions. Following the approach of our earlier papers [24, 34], we consider in Sect. 2 the leading-particle effect as a possible source of a classical pion field in the impact-parameter space. It will be related to the semiclassical solutions of the equation of motion of the nonlinear σ-model coupled to quark degrees of freedom. In order to faciliate the analysis of charged-neutral pion correlations, we also derive the corresponding pion-generating function. The quantum version of the nonlinear σ-model coupled to quarks is discussed in Sect. 3, and shown to lead to a coherent state description of the pions emerging from the DCC. A multipion exchange potential between two quarks is derived in the configuration space and its light-cone singularity structure discussed. Results of our investigation are summarized in the Sect. 4. Our general conclusion is that within the nonlinear σ-model the large isospin fluctuations depend strongly on the value of the ρ/π ratio which fluctuate from event to event.

2. *Pion production from a classical source*

At high energies most of the pions are produced in the nearly baryon-free central region. The energy available for the hadron production is

$$E_{\rm had} = \frac{1}{2}\sqrt{s} - E_{\rm leading}\,, \tag{3}$$

which at fixed total c.m. energy $\sqrt{s}$ varies from event to event. The n-pion contribution to the s-channel unitarity can be written as an integral over the relative impact parameter b of the two incident leading particles

$$F_n(s) = \frac{1}{4s} \int \mathrm{d}^2 b \prod_{i=1}^{n} dq_i \mid T_n(s, \boldsymbol{b}; 1 \dots n) \mid^2, \tag{4}$$

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

$$\begin{aligned} F_n(s) &= s\sigma_n(s) \\ \text{and} \qquad \sigma_{\text{inel}}(s) &= \sum_{n=1}^{\infty} \sigma_n(s). \end{aligned} \tag{5}$$

If the isospin of the incoming leading particles is II_3, then the initial-state vector of the pion field is $\hat{S}(s, \boldsymbol{b}) \mid II_3\rangle$, where $\mid II_3\rangle$ denotes a vacuum state with no pions. The n-pion production amplitude is

$$\mathrm{i}\, T_n(s, \boldsymbol{b}; q_1 \dots q_n) = 2s\langle I' I'_3; q_1 \dots q_n \mid \hat{S}(s, \boldsymbol{b}) \mid II_3\rangle, \tag{6}$$

where $I' I'_3$ denotes isostate of the outgoing leading particles.

The basic assumption of the independent emission of pion-clusters, in b-space is the factorization property of the scattering amplitude. The factorization of T_n follows if the quantum field of the pion-cluster satisfies the equation of motion of the form

$$(\Box + m_c^2)\boldsymbol{\pi_c}(s, \boldsymbol{b}; x) = \boldsymbol{j_c}(s, \boldsymbol{b}; x), \tag{7}$$

for an isovector cluster and similar one for an isoscalar cluster. Here, $\boldsymbol{j_c}$ is a classical source of the cluster which decays into c = 1,2,... pions outside the region of strong interactions (the final-state interaction between pions being neglected). Clusters decaying into two or more pions simulate a short-range correlation between pions. They need not be well-defined resonances.

If the conservation of isospin is a global property of the colliding system, then $\boldsymbol{j_c}(s, \boldsymbol{b}; x)$ is of the form

$$\boldsymbol{j_c}(s, \boldsymbol{b}; x) = j_c(s, \boldsymbol{b}; x)\boldsymbol{n}, \tag{8}$$

where $\boldsymbol{n}$ is a fixed unit vector in isospace independent of x, s and $\boldsymbol{b}$. The global conservation of isospin thus introduces the long-range correlation between the emitted pions.

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

The coherent production of pion-clusters is described by the S matrix

$$\hat{S}(s, \boldsymbol{b}) = \int \mathrm{d}^2 \boldsymbol{n} \mid \boldsymbol{n}\rangle \hat{D}(s, \boldsymbol{b}) \langle \boldsymbol{n} \mid, \tag{9}$$

where $|\, \boldsymbol{n} \rangle$ represents the isospin-state vector of the two-leading-particle system. The quantity $\hat{D}(s, \boldsymbol{b})$ is the unitary coherent-state displacement operator defined in our case as

$$D(s, \boldsymbol{b}) = \exp[a^{\dagger}(s, \boldsymbol{b}) - a(s, \boldsymbol{b})]\,, \tag{10}$$

where

$$a^{\dagger}(s, \boldsymbol{b}) = \sum_{c} \int \mathrm{d}q [J_c(s, \boldsymbol{b}; q) \boldsymbol{n} \boldsymbol{a_c}^{\dagger}(q) + J'_c(s, \boldsymbol{b}; q) a^{\dagger}_c(q)], \tag{11}$$

and $\boldsymbol{a_c}^{\dagger}(q)$ and $a^{\dagger}_c(q)$ are the creation operators of a cluster of type c, respectively.

The isospin (I', I'_3) of the outgoing leading particle system varies from event to event. It is produced with the probability ω_{I', I'_3}, and we can sum over all $(I' I'_3)$ to obtain the probability distribution of producing $n_+ \pi^+$, $n_- \pi^-$, and $n_0 \pi^0$ from a given isospin state

$$P_{II_3}(n_+ n_- n_0) N_{II_3} = \sum_{I' I'_3} \omega_{I', I'_3} \int \mathrm{d}^2 b \, \mathrm{d}q_1 \mathrm{d}q_2 \ldots \mathrm{d}q_n \mid \langle I' I'_3 n_+ n_- n_0 \mid \hat{S}(s, \boldsymbol{b}) \mid I I_3 \rangle \mid^2, \tag{12}$$

where $n = n_+ + n_- + n_0$ and N_{II_3} is the corresponding normalization factor determined by $\sum_{\{n\}} P_{II_3}(n_+, n_-, n_0) = 1$.

This is our basic relation for calculating various pion-multiplicity distributions, pion multiplicities, and pion correlations between definite charge combinations. In general, the probability $P_{II_3}(n_+ n_- n_0)$ would depend on $(I' I'_3)$ dynamically. The final-leading-particles tend to favour the $(I', I'_3) \approx (I, I_3)$ case. However, if the leading particles are colliding nuclei, an almost equal probability for various (I', I'_3) seems reasonable approximation owing to the large number of possible leading isobars in the final state. Then we can sum over all $(I' I'_3)$ using the group theory alone.

Recent studies of heavy-ion collisions at the partonic level [26] argue that the central region is mainly dominated by gluon jets. The valence quarks of the incoming particles which escape from the interaction region form the outgoing leading particle system. Since gluon's isospin is zero, it is very likely that 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 the extremely high collision energies.

Let us now assume that pions are produced both singly and through isovector clusters of the ρ type [24]. In this case, the most appropriate tool for studying various pion correlations is the generating function $G_{II_3}(z, n_-)$

$$G_{II_3}(z, n_-) = \sum_{n_0, n_+} P_{II_3}(n_+, n_-, n_0) z^{n_0}, \tag{13}$$

from which we can calculate, for example

$$\begin{aligned}\langle n_0\rangle_{n_-} &= \frac{\mathrm{d}}{\mathrm{dz}}\ln G_{II_3}(z,n_-)\,|_{z=1} \\ &= \sum_{n_0,n_+} n_0 P_{II_3}(n_+,n_-,n_0),\end{aligned} \tag{14}$$

$$f^0_{2,n_-} = \frac{\mathrm{d}^2}{\mathrm{dz}^2}\ln G_{II_3}(z,n_-)\,|_{z=1} \tag{15}$$

$$= \langle n_0(n_0-1)\rangle_{n_-} - \langle n_0\rangle^2_{n_-}, \tag{16}$$

and

$$\begin{aligned}P_{II_3}(n_0) &= \frac{1}{n_0!}\frac{\mathrm{d}^{n_0}}{\mathrm{d}z^{n_0}}\sum_{n_-} G_{II_3}(z,n_-)\,|_{z=0} \\ &= \sum_{n_-,n_+} P_{II_3}(n_+,n_-,n_0)\,.\end{aligned} \tag{17}$$

The form of this generating function is

$$G_{II_3}(z,n_-) = (I+\frac{1}{2})\frac{(I-I_3)!}{(I+I_3)!}\int_{-1}^{1}\mathrm{d}x\,\mid P_I^{I_3}(x)\mid^2\,\frac{A(z,x)^{n_0}}{n_0!}e^{-B(z,x)}, \tag{18}$$

where

$$2A(z,x) = (1-x^2)\overline{n}_\pi + z(1-x^2)\overline{n}_\rho + 2x^2\overline{n}_\rho \tag{19}$$

and

$$2B(z,x) = \overline{n}_\pi(1+x^2-2zx^2) + \overline{n}_\rho(2-z(1-x^2)). \tag{20}$$

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 associate Legendre polinomial. Note that $A(1,x) = B(1,x)$.

The total number of emitted pions is

$$\overline{n} = \overline{n}_\pi + 2\overline{n}_\rho. \tag{21}$$

The behaviour of $P(n_0) \equiv P_{00}(n_0)$ for $\overline{n} = 50$ and different combinations of $(\overline{n}_\pi, \overline{n}_\rho)$ is shown in Fig. 1.

The behaviour of the n_0-dispersion

$$D(n_0)^2_{n_-} = f^0_{2,n_-} + \langle n_0\rangle_{n_-} \tag{22}$$

for a given number of negative pions and different pairs of $(\overline{n}_\pi, \overline{n}_\rho)$ in the case of $I = I_3 = 1$ is shown in Fig. 2. Note that f^0_{2,n_-} is a sensitive quantity of the pairing properties of the pions.

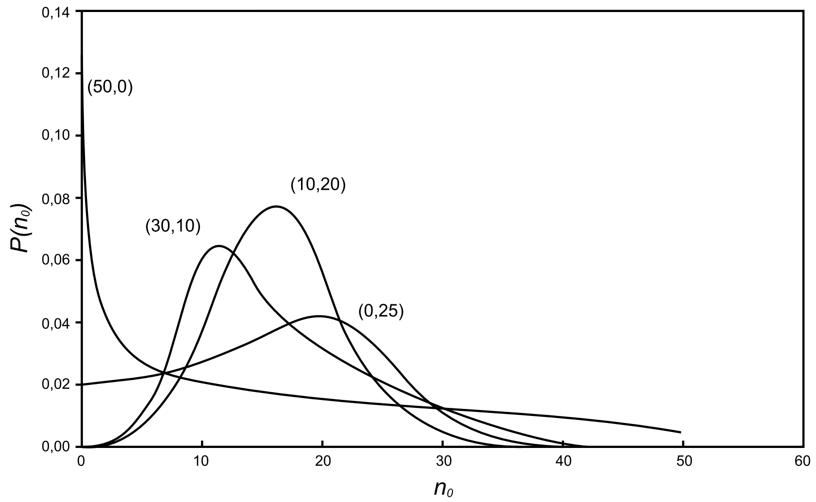

Fig. 1. The curves represent $P(n_0)$ for different combinations of $(\overline{n}_\pi, \overline{n}_\rho)$, the average number of singly produced pions and the average number of ρ-type clusters of pions, respectively.

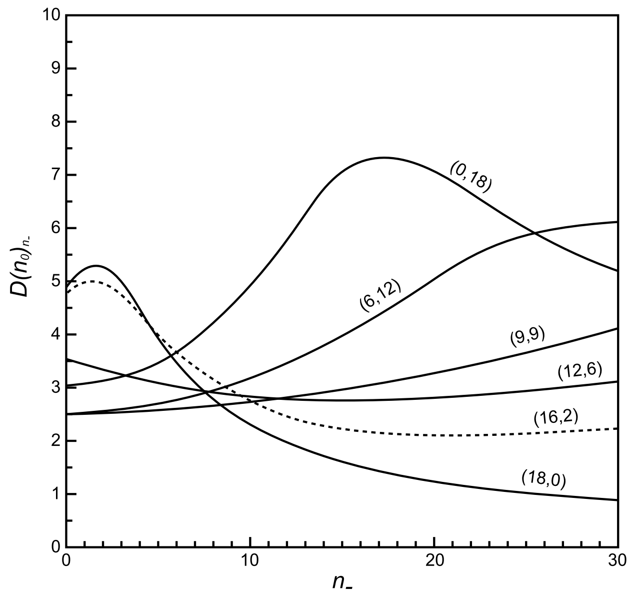

Fig. 2. Dispersion of two neutral pions for a given number of negative pions in pp-*collisions ($I = I_3 = 1$). The curves represent different combinations of $(\overline{n}_\pi, \overline{n}_\rho)$.*

We see that DCC-type behaviour is expected only for $\overline{n}_\pi \neq 0$ and $\overline{n}_\rho = 0$, that is in events without ρ-resonances. Recent estimate of the ratio of ρ– mesons to pions, at accelerator energies, is $\overline{n}_\rho = 0.10\overline{n}_\pi$ [6].

3. *Isospin fluctuation in a quantum nonlinear sigma model*

In this section, we establish the relationship between the quantum nonlinear sigma model coupled to quarks and our coherent state eikonal model description of the pion production.

As is well known, the Lagrangian for QCD with two light up and down quarks has an approximate global $SU(2)_\mathrm{L} \times SU(2)_\mathrm{R}$ symmetry, which at low temperatures, is spontaneously broken to $SU(2)_\mathrm{V}$ by a nonzero value of the quark condensate $\langle \bar{q}_\mathrm{L} q_\mathrm{R} \rangle$, which is regarded as an order parameter of the system. This order parameter can be represented as a four-component vector $\phi \equiv (\sigma, \boldsymbol{\pi})$ built from the quark densities. The chiral symmetry then corresponds to $O(4)$ rotations in internal space.

The true vacuum of the theory is defined as $\langle \phi \rangle = (\langle \sigma \rangle, \mathbf{0})$, with $\langle \sigma \rangle \neq 0$. In QCD, the spontaneous symmetry breakdown leads to nearly massless Goldstone bosons (the pions) and gives the constituent-quark mass. At low energies and large distances (momentum scale smaller than 1 GeV), the dynamics of QCD is described by an effective Lagrangian containing the $\sigma, \boldsymbol{\pi}$ fields and constituent quarks.

In the DCC dynamics, we distinguish three stages: formation, evolution and decay stage. In the conventional approach [19], one starts with a chirally symmetric phase at $T > T_\mathrm{c}$ and DCC formation happens as T, due to a rapid expansion or cooling, drops below T_c spontaneously breaking the chiral symmetry.

The evolutionary stage of the DCC is usually described by the classical chiral dynamics based on the σ-model, mostly the linear σ-model. We consider the nonlinear σ-model coupled to quarks at zero temperature [35, 36] which is expected to describe the late stage of the DCC evolution.

The Lagrangian for the nonlinear σ-model coupled to quarks is

$$L = \frac{f_\pi^2}{4} Tr(\partial_\mu U^\dagger \partial^\mu U) + \overline{q}(\mathrm{i}\,\gamma\partial)q - g f_\pi \overline{q} U q \,, \tag{23}$$

where

$$U = \exp(\mathrm{i}\,\gamma_5 \frac{\boldsymbol{\pi}\cdot\boldsymbol{\tau}}{f_\pi})\,. \tag{24}$$

We shall parametrize the pion field in the following form

$$\boldsymbol{\pi}(x) = f_\pi \boldsymbol{n}(x)\,\theta(x)\,, \tag{25}$$

where $\boldsymbol{n}(x)$ is a unit vector which determines the local isospin orientation of the pion field, obeying $\boldsymbol{n}(x)^2 = 1$.

The Euler-Lagrange equations of motion for θ and $\boldsymbol{n}$ are

$$\Box\,\theta - \sin\theta\cos\theta\partial_\mu \boldsymbol{n}\cdot\partial^\mu \boldsymbol{n} = -\mathrm{i}\,\frac{m_Q}{f_\pi^2}\boldsymbol{n}\cdot(\bar{Q}\boldsymbol{\tau}\gamma_5 Q)\,, \tag{26}$$

$$\partial_\mu(\sin^2\theta \boldsymbol{n}\times\partial^\mu \boldsymbol{n}) = -\mathrm{i}\,\frac{m_Q}{f_\pi^2}\boldsymbol{n}\times(\bar{Q}\boldsymbol{\tau}\gamma_5 Q)\sin\theta\,, \tag{27}$$

where $Q(x)$ denotes the constituent quark field defined by

$$Q(x) = \exp\left(\mathrm{i}\,\gamma_5\frac{\boldsymbol{\pi}\cdot\boldsymbol{\tau}}{2f_\pi}\right)q(x), \tag{28}$$

and $m_\mathrm{Q} = gf_\pi$ is the constituent quark mass.

We treat

$$-\mathrm{i}\,\frac{m_\mathrm{Q}}{f_\pi}\bar{Q}(x)\boldsymbol{\tau}\gamma_5 Q(x) = \boldsymbol{j}(x) \tag{29}$$

as a given classical external source of pions and identify it with the source function $\boldsymbol{j}_{(c=1)}(s,\boldsymbol{b};x)$ in our eikonal model (7). For the class of solutions that can be rotated into a uniform one, $\boldsymbol{n}(x) = \boldsymbol{n}$, known as the Anselm-class of solutions [16], the solutions with constant $\boldsymbol{n}$ can be realized if the source points to a certain fixed direction $\boldsymbol{n}$ in the isospace

$$\boldsymbol{j}(x) = j(x)\boldsymbol{n}\,. \tag{30}$$

Then the equation of motion for the pion field reduces to

$$\Box\,\theta(x) = j(x) \tag{31}$$

with

$$\boldsymbol{\pi}(x) = f_\pi\theta(x)\boldsymbol{n}\,. \tag{32}$$

This relationship offers the possibility to study the importance of various quark sources in the DCC formation.

In quantum chiral field theory of the nonlinear σ-model, the role of a strong coupling of qπ-interaction has the ratio m_Q/f_π. Since the Lagrangian of the nonlinear σ-model is nonpolynomial, a suitable renormalization procedure and operator normal ordering should be formulated [37, 38]. Let

$$\Phi =: \exp\left(\mathrm{i}\gamma_5\frac{\boldsymbol{\pi\tau}}{f_\pi}\right) - 1\,: \tag{33}$$

denote the chiral super field of the pion. The Lagrangian describing the interaction of this super chiral field with quarks is now

$$L = \frac{f_\pi^2}{4}\mathrm{Tr}(\partial_\mu\Phi\partial^\mu\Phi) + \bar{q}(\mathrm{i}\,\gamma\partial - m_\mathrm{Q})q - \frac{m_\mathrm{Q}}{f_\pi}\bar{q}\Phi q. \tag{34}$$

If the Φ field is expanded in terms of pion fields

$$\Phi = \frac{\mathrm{i}}{\mathrm{f}_\pi}\gamma_5 \boldsymbol{\pi}\cdot\boldsymbol{\tau} - \frac{1}{2f_\pi^2}\boldsymbol{\pi}\cdot\boldsymbol{\pi} + \ldots\,, \tag{35}$$

the first term will generate one-pion exchange interaction between quarks. The main contribution of the second term will come from the two-pion exchange which can be simulated by a sigma-meson exchange. The nonrelativistic reduction of the above Lagrangian gives an interaction between quarks which is mediated by the exchange of an arbitrary number of pions grouped in different Goldstone-type bosons.

However, for studying multipion production on quarks, it is necessary to find the chiral superpropagator of the $\Phi(x)$ field

$$\begin{aligned} \triangle_\Phi(x) &= \langle T(\Phi(x)\Phi(0))\rangle \\ &= \langle T(:\exp\left(\mathrm{i}\,\gamma_5\frac{\boldsymbol{\pi}(x)\boldsymbol{\tau}}{f_\pi}\right)::\exp\left(\mathrm{i}\,\gamma_5\frac{\boldsymbol{\pi}(0)\boldsymbol{\tau}}{f_\pi}\right):\rangle - 1\,. \end{aligned} \tag{36}$$

Its explicit form in x-space, after a number of algebraic manipulations, is

$$\triangle_\Phi(x) = 1\otimes 1F(x) + \frac{1}{3}(\gamma_5\boldsymbol{\tau})\otimes(\gamma_5\boldsymbol{\tau})G(x)\,, \tag{37}$$

where

$$F(x) = \mathrm{ch}\left(\frac{\triangle(x)}{f_\pi^2}\right) + \frac{\triangle(x)}{f_\pi^2}\,\mathrm{sh}\left(\frac{\triangle(x)}{f_\pi^2}\right) - 1 \tag{38}$$

$$\text{and}\quad G(x) = 2\,\mathrm{sh}\left(\frac{\triangle(x)}{f_\pi^2}\right) + \frac{\triangle(x)}{f_\pi^2}\,\mathrm{ch}\left(\frac{\triangle(x)}{f_\pi^2}\right)\,,$$

where the massless pion propagator

$$\triangle(x) = \frac{1}{4\pi^2}\frac{1}{x^2 - \mathrm{i}\epsilon}$$

is singular on the light-cone. Due to the singular character of the $\triangle(x)$, the superpropagator $\triangle_\Phi(x)$, without further restrictions, contains an infinite number of arbitrary parameters which can be expressed as a δ-ambiguity

$$\delta(\triangle_\Phi(x)) = \sum_n a_n \Box^n \delta^{(4)}(x)\,, \tag{39}$$

with a_n real and $\sum a_n z^n$ an entire function of order $\leq 1/2$. A definite choice of these parameters can be made [37], so that $\triangle_\Phi(x)$ has no real singularities of the type $\Box^n\delta^{(4)}(x)$. The trick is to work in the Euclidean space-time, where no $\mathrm{i}\epsilon$ appears in the propagator $\triangle(x)$.

This superpropagator can be used to study a contribution to the quark-quark scattering amplitude due to the exchange of a single chiral superfield. The multipion exchange potential between two quarks is related to the Fourier transform of $\triangle_\Phi(x)$ in the following way

$$\bar{u}(p_1')\bar{u}(p_2')\tilde{\triangle}_\Phi(q)u(p_1)u(p_2) = \omega_1'\omega_2'\tilde{V}_\Phi(q)\omega_1\omega_2\,, \tag{40}$$

where $q = p_1' - p_1 = p_2' - p_2$ and

$$u(p) = \sqrt{\frac{\epsilon_p + m}{\epsilon_p}}\begin{pmatrix}\omega\\ \omega'\end{pmatrix}, \quad \omega' = \frac{\boldsymbol{\sigma p}}{\epsilon_p + m}\omega$$

with $\omega^*\omega = 1$. For space-like momentum-transfer variable $(q, q^2 \leq 0)$, the behaviour of the Fourier transform of $\mathrm{ch}(\triangle(x)/f_\pi^2)$, for example, is $\mathrm{ch}(\mid q/f_\pi \mid^{2/3})$ for $q \to \infty$.

According to the usual definition of the potential, as a nonrelativistic limit of the Fourier-transformed exchange diagram in the field theory

$$V_\Phi(\boldsymbol{r}) = (2\pi)^{-3}\int \mathrm{d}^3q\; e^{\mathrm{i}\,\boldsymbol{qr}}\tilde{V}_\Phi(0,\boldsymbol{q})\,, \tag{41}$$

we distinguish two contributions: one coming from the $F(x)$-part of the superpropagator, $\triangle_\Phi(x)$, and the other one coming from the $G(x)$-part, respectively. F generates central and spin-orbit forces between two quarks while G generates spin-spin and tensor forces.

Over the past years, the one and two boson-exchange potentials (BEP) in NN-dynamics and their extension to Goldstone-boson-exchange (GBE) dynamics have already been successfully implemented in a chiral quark model for baryons [39]. Our approach, using the chiral superfield, seems to be a step forward in the phenomenological description of the quark-quark interaction.

In the limit of soft chiral pion bremsstralung, which was studied long time ago [40 – 43], every incoming and outgoing quark line was replaced by

$$q(x) \longrightarrow Q(x) = \exp(\mathrm{i}\,\gamma_5\frac{\boldsymbol{\pi\tau}}{2f_\pi})q(x). \tag{42}$$

In this case, the distribution of neutral pions from a single quark line follows the form $P(n_0 \mid n) \sim \sqrt{1/nn_0}$ [41], which is a typical behaviour for coherent pion production without invoking the notion of DCC formation.

4. Conclusion

The results of the present analysis have shown that the experimental observation of DCC is strongly affected by the ρ/π production ratio and the soft chiral-pion bremsstralung. In particular, we have found that:

- Within the framework of a unitary eikonal model with factorization, energy conservation and global conservation of isospin the DCC-type fluctuation of the neutral pion fraction (f) could be obtained if the ρ/π ratio is small in an event, at least in a restricted phase-space domain;
- The DCC effect, if observed would also depend on isospin of the initial-leading-particle system;
- The coherent production of ρ-type clusters of pions suppresses the DCC-type fluctuations;
- The factorization property of the scattering amplitude in the impact-parameter space of the initial-leading-particle system can be related to the isospin-uniform solutions of the quantum nonlinear σ-model coupled to quarks; relations (7) and (8), and (29) – (31);
- The multipion exchange potential between two quarks in the configuration space can be derived. If its light-cone singularities are properly treated, it can be used to study multipion production processes in terms of the ratio m_{Q}/f_π;
- The soft chiral pion bremsstralung also leads to anomalously large ratio of neutral to charged pions. Therefore, the large n behaviour of $P(n_0 \mid n) \sim \sqrt{1/nn_0}$ is not a definite signature of DCC formation.

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DEZORIJETIRANI KIRALNI KONDENZAT I FLUKTUACIJE NABIJENIH I NEUTRALNIH ČESTICA U SUDARIMA TEŠKIH IONA

Proučavali smo mogućnost velikih fluktuacija naboja i izospina tipa DCC (DCC = dezorijentirani kiralni kondenzat) u visokoenergijskim sudarima teških iona u okviru unitarnog eikonalnog modela. Primijenili smo svojstvo faktorizacije amplitude raspršenja u prostoru parametra sudara vodećih čestica za proučavanje poluklasične tvorbe piona u centralnom području. Funkcija klasičnog izvora pionskog polja u prostoru parametra sudara može se povezati s poluklasičnim rješenjima jednadžbe gibanja u nelinearnom sigma-modelu vezanom na kvarkovske stupnjeve slobode. U radu se izvodi izraz za potencijal višepionske izmjene između dva kvarka. Pokazuje se da se u limesu kočnog zračenja mekanih kiralnih piona također pojavljuju anomalno velike fluktuacije omjera neutralnih i nabijenih piona neovisno o pojavi DCC-a. Također se pokazuje da bi fluktuacije tipa DCC trebale biti potisnute u slučaju kada se pioni proizvode uglavnom putem raspada ρ-mezona.