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Critical events and intermittency in nuclear collisions

N.G. Antoniou [1], F.K. Diakonos, C.N. Ktorides, M. Lahanas

Department of Physics, University of Athens, GR-15771 Athens, Greece

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Abstract

The structure of events in relativistic heavy ion collisions (RHIC), associated with chiral QCD phase transition, is investigated. In particular the density fluctuations of pions, emitted from the excited vacuum (at $T \approx T_c$), are studied and classified in terms of intermittency patterns in rapidity space. For this purpose a Monte Carlo simulation of critical events in the central region is presented, on the basis of the $O(4)$ theory of chiral phase transition. The universal character of these events is revealed and the predictions of the theory are discussed in connection with event by event measurements in current and future experiments with relativistic heavy ions.

1. Introduction

A prime, general indicator of a quark-hadron phase transition in central A + A collisions is the appearance of critical dynamical fluctuations in individual events [1]. In particular, one expects a universal pattern of multiplicity density fluctuations at midrapidity, whose origin can be ascribed to the critical behaviour of the underlying theory of strong interactions (chiral QCD at $T = T_c$). The interrelation between the fundamental parameters of the theory (critical exponents) and the measurable indices of these events (intermittency exponents) becomes of primary importance in our effort to understand the formation of a new phase of matter (QGP) in relativistic heavy ion collisions (RHIC).

In this work a class of such critical events is derived, directly from the universality class of chiral QCD phase transition in the limit of infinite strange quark mass ($m_u \approx m_d \approx 0, m_s = \infty$). In this case the system undergoes a second-order phase transition, belonging to the $O(4)$ universality class [2] and the corresponding order parameter is specified by the sigma (σ) and pion $(\boldsymbol{\pi})$ condensates, $\phi_\alpha = (\boldsymbol{\pi}, \sigma)$. Our approach is applicable even if the mass of the strange quark is finite, provided that it remains larger than a critical value m^* corresponding to a tricritical point of the system [2]. If $m_s < m^*$, the system undergoes a first order transition and a different pattern of fluctuations is expected to appear near the critical temperature T_c [3]. A comprehensive study of the critical fluctuations, taking into account the role of the strange quark mass, will be presented elsewhere.

[1] E-mail address: nantonio@atlas.uoa.gr.

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Our aim, in this Letter, is to show that the experimental measurement of a class of events in A + A collisions, displaying, in the distribution of pions, the characteristic fluctuations of the $O(4)$ critical theory, could provide us with a strong evidence for the formation of QGP in the collision and the occurence of a second-order chiral phase transition in the process of hadronization. For this purpose we adopt the standard description of a longitudinal expansion after the collision (inside-outside dynamics) and focus our attention on the 1d projection onto rapidity space (ξ-space). We assume that local thermal equilibrium is reached rather soon, so at $T = T_c$ the onset of a static phase transition of second order (chiral) drives the system smoothly from the quark phase (massless u,d quarks) to a phase of primordial hadrons (massless π, σ) with a macroscopic density given by the classical order parameter $\phi^2 = \boldsymbol{\pi}^2 + \sigma^2$ [4]. Our approach is complementary to the proposal for DCC production occuring in a chiral phase transition far from equilibrium [5].

2. Intermittency at $T = T_c$

The $O(4)$ effective action (at $T = T_c$) for a 3d hadronization source localized at $\xi = \xi_i$ and a proper time τ, is written, in terms of the classical fields ϕ_α, as follows:

$$\Gamma_c\left[\hat{\phi}\right] = C_A \int dz d^2\boldsymbol{\chi}_\perp \left[\frac{1}{2}\left(\frac{\partial\hat{\phi}}{\partial z}\right)^2 + \tfrac{1}{2}\left(\nabla_\perp \hat{\phi}\right)^2 + \frac{3C_A^2}{4}\left(\hat{\phi}^2\right)^\kappa \right],$$

$$z = \sinh(\xi - \xi_i), \quad \boldsymbol{\chi}_\perp = \frac{\boldsymbol{x}_\perp}{\tau}, \quad C_A = \frac{\tau}{\beta_c},$$

$$\hat{\phi}_\alpha = \beta_c \phi_\alpha,$$

$$\hat{\phi}^2 = \sum_{\alpha=1}^{4} \hat{\phi}_\alpha^2, \quad \left(\partial\hat{\phi}\right)^2 = \sum_{\alpha=1}^{4}\left(\partial\hat{\phi}_\alpha\right)^2. \qquad (1)$$

The effective potential $U_c \sim (\hat{\phi}^2)^\kappa$ in Eq. (1) is the scaling solution at $T = T_c$ of the $O(4)$ theory [6] given in terms of the anomalous dimension $\eta \approx 0.034$ ($\kappa = \frac{3}{1+\eta}$). Projecting out onto rapidity space we may neglect the details in transverse directions assuming $\nabla_\perp \phi_\alpha \approx 0$, whereas for sufficiently small size Δ of the source in ξ-space, we may write $dz \approx d\xi$ in Eq. (1). Hence, the effective action of the 1d system in rapidity space is written as follows:

$$\Gamma_c\left[\hat{\phi}\right] = \frac{\pi R_\perp^2}{\tau\beta_c} \int_\Delta d\xi \left[\tfrac{1}{2}\left(\frac{\partial\hat{\phi}}{\partial\xi}\right)^2 + \frac{3C_A^2}{4}\left(\hat{\phi}^2\right)^\kappa \right]. \qquad (2)$$

Although the form (2) is strictly valid for $\Delta \ll 1$, its analytic continuation may accommodate sources at $\xi = \xi_i$ of any size Δ available in rapidity space. It defines the free energy at $T = T_c$, associated with the formation in ξ-space of chiral condensates $(\boldsymbol{\pi}, \sigma)$ from a hadronization source located at $\xi = \xi_i$. The Statistical Mechanics of this process is described by the partition function $Z = \int[\mathscr{D}\hat{\phi}_\alpha]e^{-\Gamma_c[\hat{\phi}]}$ and for heavy nuclei ($\frac{\pi R_\perp^2}{\tau\beta_c} \gg 1$) one should look for saddle-point configurations ($R_\perp$ is the transverse radius of the 3d cylindrical source, at $T = T_c$). For a given orientation of the vacuum in the $O(4)$ space, $\hat{\phi}_\alpha(\xi) = \lambda_\alpha \phi(\xi)$, $\sum_{\alpha=1}^4 \lambda_\alpha^2 = 1$, the solutions of the Euler-Lagrange equation, $\ddot{\phi} - \frac{3\kappa C_A^2}{2}\phi^{2\kappa-1} = 0$, give the saddle-point configurations which are classified according to the energy E of the corresponding classical trajectories. In what follows the trajectories with $E \neq 0$ are neglected since they contribute to the partition function with a suppression factor of the form $Z(E) \sim \exp(-\frac{V_s}{\beta_c\tau}|E|)$, where $V_s = \pi R_\perp^2 \Delta$ is the volume of the 3d source in the cylindrical space $(\xi, \boldsymbol{x}_\perp)$. The dominant configurations ($E = 0$) and their contribution to the free energy $\Gamma_c[\phi]$ are written as follows:

$$\phi(\xi) = \left[\frac{1}{C_A(\kappa-1)}\sqrt{\frac{2}{3}}\right]^{\frac{1}{\kappa-1}} |\xi - \xi_o|^{-\frac{1}{\kappa-1}},$$

$$\Gamma_c(\Delta\xi; \xi_o) = \frac{3\pi C_A R_\perp^2}{2\beta_c^2} \int_{\Delta\xi} [\phi(\xi)]^{2\kappa} d\xi. \qquad (3)$$

The location of the singularity ($\xi = \xi_o$) is arbitrary but, for a given ξ_o, the solution (3) gives the saddle-point configurations (with different supports) only for systems localized in domains $\frac{\Delta\xi}{2} < |\xi_o - \xi_i|$ and emitted from the source at $\xi = \xi_i$. The restriction on the size $\Delta\xi$ follows from the fact that the singularity must lie outside the domain of the system since otherwise it would give an infinite contribution to the effective action. Considering now the whole family of solutions (3) one might be able to reveal the critical fluctuations of the system in a wide range of scales ($\Delta\xi$). For this purpose we introduce, in what follows, subdivisions of the size of the system $\Delta\xi = \frac{\Delta}{m}$ ($m \geq 1$) in order to specify the support of the saddle-point configurations contributing to the partition function. A further restriction on the configurations follows from the observation that only distant singularities ($\xi_o \gg \xi_i + \frac{\Delta\xi}{2}$) give a significant contribution to the partition function. In fact the free energy corresponding to the solution (3), $\Gamma_c(\Delta\xi;\xi_o) \sim (\xi_o - \xi_i - \frac{\Delta\xi}{2})^{-\frac{\kappa+1}{\kappa-1}}$, increases rapidly when ξ_o approaches the end point $\xi_i + \frac{\Delta\xi}{2}$. Hence, one may safely use constant configurations $\phi \approx [\frac{\sqrt{2/3}}{\xi_o C_A(\kappa-1)}]^{1/(\kappa-1)}$ in domains $\Delta\xi = \frac{\Delta}{m}$, in order to compute the partition function: $Z = \sum_{m,\xi_o} e^{-\Gamma_c(m;\xi_o)}$. Introducing the total multiplicity n_s of hadrons as a suitable random variable in describing the critical fluctuations of the system, we recall that ϕ^2 is proportional to the rapidity density $\rho(\xi,\xi_i)$ of hadrons emitted from the source at ξ_i as a mixture of pions and sigmas according to the orientation $\{\lambda_\alpha\}$ of the vacuum in the internal 4-dimensional space: $\phi^2 = \frac{1}{\pi}(\frac{\beta_c}{R_\perp})^2\rho(\xi,\xi_i)$ [4]. With this interpretation the saddle-point, constant field ϕ, can be expressed in terms of the random variables (n_s,m), which belong to the grand canonical ensemble of the system [4]. After integrating over the orientations $\{\lambda_\alpha\}$, the canonical partition function at $T = T_c$ takes the form:

$$Z(n_s,\Delta,T_c) = \pi^2 n_s \sum_{m\geq 1} \exp\left[-\left(\frac{\delta_o}{\Delta}\right)^{\kappa-1} n_s^\kappa m^{\kappa-1}\right],$$

$$\delta_o = \frac{\beta_c^2}{\pi R_\perp^2}\left(\frac{3C_A}{4}\right)^{\frac{1}{\kappa-1}}. \tag{4}$$

The partition function (4) gets contribution from the dominant saddle point configurations associated with the $O(4)$ effective potential and, therefore, provides us with a satisfactory representation of the 1d projection of the critical system. Using Eq. (4) one may extract the average multiplicity $\langle n_s\rangle$ and the corresponding density-density correlation $\rho(\xi,\xi_i) \equiv \langle\rho(\xi)\rho(\xi_i)\rangle$ associated with the source at $\xi = \xi_i$:

$$\langle n_s\rangle = \frac{(\kappa-2)\,\Gamma\left(\frac{3}{\kappa}\right)}{(2\kappa-3)\,\Gamma\left(\frac{2}{\kappa}\right)}\left(\frac{\Delta}{\delta_o}\right)^{\frac{\kappa-1}{\kappa}},$$

$$\rho(\xi,\xi_i) = \frac{(\kappa-1)(\kappa-2)\,\Gamma\left(\frac{3}{\kappa}\right)}{2\kappa(2\kappa-3)\,\Gamma\left(\frac{2}{\kappa}\right)} \times\left(\frac{2}{\delta_o}\right)^{\frac{\kappa-1}{\kappa}}|\xi-\xi_i|^{-\frac{1}{\kappa}}. \tag{5}$$

The power-laws (5), corresponding to a fractal dimension $d_F^{(1)} = \frac{\kappa-1}{\kappa}$ in rapidity space, are valid for $\delta_o \ll |\xi-\xi_i| \leq \Delta$ and have a universal character. As a result we claim that a characteristic process in relativistic nuclear collisions, associated with chiral QCD phase transition, is the emission of a universal intermittent spectrum of particles, from any local hadronization source along the space-time hyperbola $\tau =$ const. The intermittency exponents, $f_q = \frac{(q-1)(1+\eta)}{3}$ ($q = 2,3,..$), are given in terms of the anomalous dimension η of the $O(4)$ theory and a nonuniversal scale δ_o, below which ($\delta\xi \leq \delta_o$) intermittency breaks down, is fixed by the finite size of the system ($\delta_o \sim \frac{\beta_c^2}{R_\perp^2}$).

3. Critical events

The self-similar profile of the hadronization sources may lead to a particular class of events in A + A collisions by considering a collection of such sources distributed at random positions, $\xi_1,\ldots,\xi_N$, according to their free energy $\Gamma_c[\langle\phi^2\rangle_s]$, given by Eq. (3) with $\langle\phi^2\rangle_s = \frac{\rho(\xi,\xi_i)}{\beta_c^{-2}\pi R_\perp^2}$. The partition function for N nonoverlapping sources is written:

$$Z_N = \sum_{\xi_1,\ldots,\xi_N}\prod_{i=1}^{N}\exp\left[-\delta_o^{\kappa-1}\left(\frac{\pi R_\perp^2}{\beta_c^2}\right)^\kappa \times\int_{\Delta\xi_i} d\xi\left(\langle\phi^2\rangle_s\right)^\kappa\right], \tag{6}$$

where the domain $\Delta\xi_i$ is fixed by the location of the sources (or the end points) in the neighbourhood of ξ_i. For $2 \le i \le N-1$ we have $\Delta\xi_i = \frac{\xi_{i+1}-\xi_{i-1}}{2}$ whereas $\Delta\xi_1 = \frac{\Delta}{2} + \frac{\xi_1+\xi_2}{2}$, $\Delta\xi_N = \frac{\Delta}{2} - \frac{\xi_N+\xi_{N-1}}{2}$. Due to the smallness of the $O(4)$ anomalous dimension ($\eta \approx 0.034$) we take in what follows, $\kappa \approx 3$. Using Eqs. (5) and (6) we finally obtain:

$$Z_N = \int_\Delta d\xi_1 \dots d\xi_N \exp\left[-\frac{1}{G}\sum_{i=2}^{N} \upsilon(\xi_i - \xi_{i-1}) - \frac{1}{2G}\left(w\left(\frac{\Delta}{2}+\xi_1\right) + w\left(\frac{\Delta}{2}-\xi_N\right)\right)\right],$$
$$-\frac{\Delta}{2} \le \xi_1 \le \dots \le \xi_N \le \frac{\Delta}{2},$$
$$G = \left[\tfrac{9}{2}\Gamma\left(\tfrac{2}{3}\right)\right]^3, \tag{7}$$

where $\upsilon(\xi) = \ln|\frac{\xi}{\delta_o}|$, $w(\xi) = \ln|\frac{2\xi}{\delta_o}|$.

The partition function (7) describes a 1-d Feynman-Wilson gas of hadronization sources with nearest neighbour ''interaction'' of the form $\upsilon(\xi_i - \xi_{i-1})$ and a similar ''potential'' $w(\frac{\Delta}{2} \pm \xi)$ giving the interaction of the chain with the walls (end points) $\xi = \pm\frac{\Delta}{2}$. The system behaves almost as an ideal gas ($G \gg 1$) with a uniform density $\frac{dN}{d\xi} \approx 1$ (one source per unit of rapidity) and a Poisson multiplicity distribution $Z_N = \frac{(\Delta)^N}{N!}$. In fact, the exact expression of the multiplicity distribution, as obtained from Eq. (7) is written:

$$P_N = \frac{Z_N}{\sum_{N=1}^{\infty} Z_N},$$

$$Z_N = \Delta^N 2^{-1/G}\left(\frac{\delta_o}{\Delta}\right)^{\frac{N}{G}} \times \prod_{i=1}^{N-1} B\left(i - \frac{2i-1}{2G}, 1-\frac{1}{G}\right) \times B\left(N - \frac{2N-1}{2G}, 1 - \frac{1}{2G}\right), \tag{8}$$

and in the limit $G \gg 1$ the beta functions appearing in (8) lead to a Poisson distribution with $\langle N \rangle \approx \Delta$. Correspondingly, the rapidity distribution for a given multiplicity N is written:

$$P(\xi_1,\dots,\xi_N) = \left(\frac{\Delta + 2\xi_1}{\delta_o}\right)^{-\frac{1}{2G}}\left(\frac{\Delta - 2\xi_N}{\delta_o}\right)^{-\frac{1}{2G}} \times \prod_{i=2}^{N}\left(\frac{\xi_i - \xi_{i-1}}{\delta_o}\right)^{-\frac{1}{G}},$$
$$-\frac{\Delta}{2} \le \xi_1 \le \dots \le \xi_N \le \frac{\Delta}{2}, \tag{9}$$

and, when combined with the self-similar profile (5) of each source, it may lead to a global structure in the central region with particular properties and characteristic patterns of critical fluctuations.

Although the effective action (1) has been taken at $T = T_c$, the resulting universal behaviour of critical fluctuations remains practically unchanged at the freeze-out temperature T_f since in the process of a second-order transition we expect that T_f remains close to T_c ($\frac{T_c - T_f}{T_c} \ll 1$). In fact one may verify this conjecture by considering a mass term $\mu^2\hat{\phi}^2$ in the free energy (2) with the standard temperature dependence $\mu^2 = \mu_o^2(\frac{T-T_c}{T_c})$. With this additional term, the partition function (4) for a subsystem of size $\delta \le \Delta$, is appropriately modified and at $T = T_f$ ($T_f < T_c$) one obtains:

$$Z(n_s,\delta,T_f) = \pi^2 n_s \sum_{m \ge 1} \exp\left[-\frac{\mu_f^2 n_s}{C_A^{(f)}} - \left(\frac{\delta_o}{\delta}\right)^{\kappa-1} n_s^\kappa m^{\kappa-1}\right], \tag{10}$$

where $\mu_f^2 = \mu_o^2(\frac{T_f - T_c}{T_c})$, $C_A^{(f)} = \frac{\tau_f}{\beta_c}$ ($\beta_f \approx \beta_c$) and τ_f is the freeze-out time-scale.

The average multiplicity of condensates radiated from the local source at $T = T_f$, is given now by a modified scaling form, valid for $\delta \gg \delta_o$:

$$\langle n_s \rangle = \left(\frac{\delta}{\delta_o}\right)^{\frac{\kappa-1}{\kappa}} F\left(\frac{|\mu_f^2|}{C_A^{(f)}}\left(\frac{\delta}{\delta_o}\right)^{\frac{\kappa-1}{\kappa}}\right), \tag{11}$$

where $F(x) = \frac{I_3(x)}{I_2(x)}$, $I_q(x) \equiv \sum_{m\ge 1}\int_0^\infty dz\, z^{\frac{q-\kappa}{\kappa}} \exp\left[-(zm^{\kappa-1} + z^{\frac{1}{\kappa}}x)\right]$ Eqs. (11) show that the deviations at $T = T_f$ from the exact power laws (5), valid at $T=T_c$, are considerably suppressed, especially for events with sufficiently long freeze-out time-scale ($C_A^{(f)} \gg 1$). In fact for $\frac{\mu_0^2}{C_A^{(f)}}\left(\frac{T_c-T_f}{T_c}\right) \ll 1$, the modified scaling form (11) reduces to the exact power laws (5) and therefore, in these events, the fractality of the system survives, to a very good approximation, even at $T=T_f$. In practice, for conventional values of the nonuniversal parameters of the thermal system: $\frac{T_c-T_f}{T_c} \approx 0.1$, $C_A^{(f)} \approx 10$, one finds a suppression of the self-similarity breaking effect by a factor of the order of 10^{-2}. Hence by ignoring intermittency breaking effects at the level of 10% at most, one may safely trust the power-laws (5) when describing the fluctuations of the system at $T=T_f$.

In order to keep contact with experimental measurements, we have also to take into account fluctua-

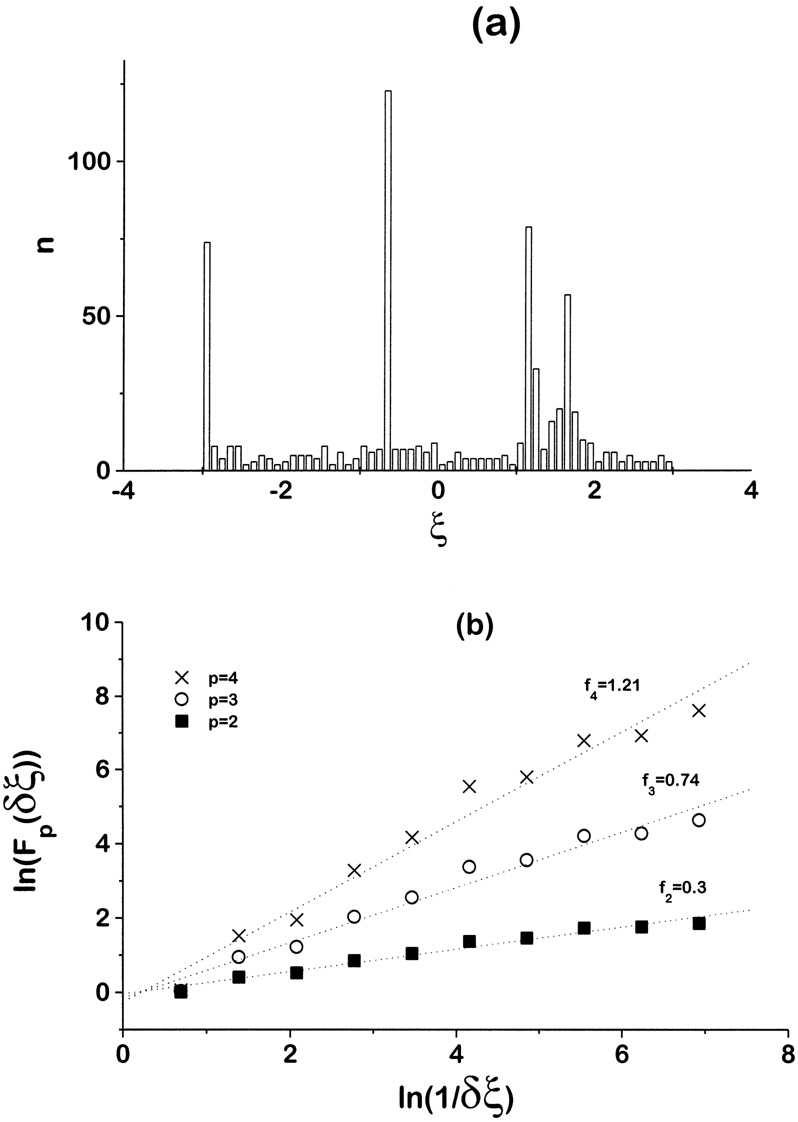

Fig. 1. (a) The rapidity spectrum of the scalar particles (π,σ) in a typical, Monte Carlo generated, critical event. The total scalar particle multiplicity is $n=687$ and the number of self-similar sources contributing to this event is $N=6$. (b) The first three factorial moments for the event presented in (a). A linear fit (dashed line) for each moment, as well as the slope of the corresponding straight line, are also shown.

tions in the transverse size of the sources $R_{\perp}^{(s)}$ expecting that on the average $\langle R_{\perp}^{(s)} \rangle \approx \beta_c$. This estimate can be derived by generalizing (6) in the 3-d cylindrical space and realizing that the corresponding 3-d sources follow also, to a very good approximation, a Poisson distribution [7]. The corresponding average multiplicity $\langle N_{\perp} \rangle$ of the sources, for a fixed rapidity, is approximately $\langle N_{\perp} \rangle \approx \left(\frac{R_{\perp}}{\beta_c}\right)^2$, suggesting again an average size for each source $\langle R_{\perp}^{(s)} \rangle \approx \beta_c$.

Taking now into account the presence of many sources, distributed in the transverse plane, the power-laws (5) are modified by changing only the nonuniversal scale-parameter in rapidity δ_o to a new effective one $\delta_o^{\rm eff} = \frac{1}{\pi}\left(\frac{\beta_c}{R_{\perp}}\right)^{\frac{2\kappa}{\kappa-1}}\left(\frac{3C_A}{4}\right)^{\frac{1}{\kappa-1}}$. It is of interest to note that, averaging over the transverse positions of the sources, the multiplicity fluctuations in the transverse plane are smeared out leading to an approximately constant density near the centre (far from the transverse boundaries). On the contrary the rapidity position of the sources is fixed on the space-time hyperbola, at the same point $\xi = \xi_i$ (independent of the transverse position), as required by our basic consideration of local thermal equilibrium. As a consequence the singularity $\rho(\xi,\xi_i) \sim |\xi - \xi_i|^{-\frac{1}{\kappa}}$ is not smoothed out by this averaging and therefore the universal characteristics of the fluctuations, described by the exponents in power-laws (5), are not affected by this modification. A detailed treatment of the critical fluctuations in the 3-d system of chiral condensates in the cylindrical geometry of relativistic heavy-ion collisions, will be presented elsewhere [7].

In our attempt to reveal the phenomenological features of the chiral QCD phase transition, a realization of the theory at the level of individual events in A + A collisions is needed. To this end, a Monte Carlo simulation (Metropolis algorithm) of critical events in rapidity space, obeying the power-laws (5) with properly corrected minimum scale δ_o and the distribution (9), has been performed. As a typical process we have chosen Pb + Pb at SPS energies ($\Delta \approx 6$) and for the scale $\delta_o^{\rm eff}$ we have taken $\delta_o^{\rm eff} \approx 10^{-4}$ ($\tau_f \approx 20$ fm, $R_{\perp} \approx 25$ fm). In Fig. 1a the rapidity spectrum of pions and sigmas in a typical event, occuring with a large probability in the Monte Carlo event generation, is presented. The number of self-similar sources contributing to this event is $N = 6$ ($\approx \Delta$) and the total multiplicity of scalars $(\sigma, \boldsymbol{\pi})$ is $n = 687$. In Fig. 1b, the first three factorial moments (scaled) $F_p(\delta\xi)$ for this event, evaluated in cells $\delta\xi$ within the central region $|\xi| \leq 2$, show a strong intermittency effect [8]. As expected, the intermittency exponents f_p follow the trend of the corresponding indices of each source ($f_p \approx \frac{p-1}{\kappa}$). Averaging over many critical events, it turns out that the rapidity distribution becomes smooth but the intermittency effect remains strong, reflecting the underlying fluctuations of individual events [7].

4. Conclusions

In conclusion, we have shown how the characteristics of the $O(4)$ universality class of chiral QCD phase transition can be converted into a universal pattern of critical fluctuations, in relativistic nuclear collisions. We have confined our treatment to the 1d projection onto rapidity space and have integrated out the details in the transverse space. A strong 1d intermittency pattern, associated with the anomalous dimension of the $O(4)$ theory, has been revealed and the details of the density fluctuations have been produced in a Monte Carlo simulation of individual events. The emerging physical picture is the emission of pions from a system of self-similar independent sources. The predicted intermittency effect, as measured by the spectrum of indices f_p, is one order of magnitude stronger than the corresponding effect observed in average events (mainly conventional), in current experiments [9]. Despite the simplifications made in our treatment, we claim that the universal characteristics of the rapidity density fluctuations, found in our approach, are sharply defined properties directly related to quark-hadron phase transition. A search for such critical events may contribute to an attempt to convert the theory of a second-order chiral QCD phase transition into a falsifiable hypothesis to be tested in relativistic nuclear collisions.

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References

[1] M. Gyulassy, Nucl. Phys. A 400 (1982) 31c; L. Van Hove, Z. Phys. C 27 (1985) 135.
[2] F. Wilczek, Int. J. Mod. Phys. A 7 (1992) 3911.
[3] N.G. Antoniou, Y.F. Contoyiannis, F.K. Diakonos, C.N. Ktorides, M. Lahanas, J. Phys. G 23 (1997) 1953.
[4] N.G. Antoniou, C.N. Ktorides, I.S. Mistakidis, F.K. Diakonos, Phys. Lett. B 353 (1995) 501; Eur. Phys. J. C, to be published.
[5] J.D. Bjorken, Int. J. of Mod. Phys. A 7 (1992) 18, 4189; K. Rajagopal, F. Wilczek, Nucl. Phys. B 379 (1993) 395; B 404 (1993) 577; J.-P. Blaizot, A. Krzywicki, Phys. Rev. D 50 (1994) 442.
[6] M. Reuter, N. Tetradis, C. Wetterich, Nucl. Phys. B 401 (1993) 567; N. Tetradis, C. Wetterich, Nucl. Phys. B 422 (1994) 541; H-Th. Elze, Nucl. Phys. A 566 (1994) 571c; N. Tetradis, private communication.
[7] N.G. Antoniou, F.K. Diakonos, C.N. Ktorides, M. Lahanas, C.G. Papadopoulos, in preparation.
[8] A. Bialas, R. Peschanski, Nucl. Phys. B 273 (1986) 703; B 308 (1988) 857.
[9] E.A. De Wolf, I.M. Dremin, W. Kittel, Phys. Rep. 270 (1996) 1.