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National Superconducting Cyclotron Laboratory

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AUREL BULGAC and LEONID FRANKFURT

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Aurel BULGAC and Leonid FRANKFURT

Department of Physics and Astronomy and National Superconducting Cyclotron Laboratory, Michigan State University, East Lansing MI 48824-1321, USA

ABSTRACT

We argue that color coherent effects lead to specific collective phenomena in the high energy heavy ion collisions with large E_t as a trigger for the centrality of the collisions. The percolation phase transition in heavy ion collisions is considered and ***methods of investigation of such a matter are discussed.***

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There are little doubts that quantum chromodynamics (QCD) is relevant for the strong interactions between hadrons. Asymptotic freedom in perturbative QCD [1] as well as more detailed predictions of pQCD have been confirmed in numerous hard processes. One of the pressing problems now is to develop a theoretical framework that will allow us to use hard processes in order to investigate softer phenomena, to search for possible phase transitions in superdense hadron matter. The aim of the present paper is to show how color coherence phenomena may be used for this purpose.

In order to convey the main idea of the paper let us consider the scattering of a sufficiently energetic composite particle h from a target T at rest. In this case different quark-gluon configurations $|n>$ in the wave function of the hadron h, with the invariant mass M_n satisfying the following condition

$$2E_h/(M_n^2 - M_h^2) \gg 2r_T \quad (1)$$

are frozen during the collision. Here r_T is the radius of the target T. Under the condition given by Rel. (1), the time necessary for a transition between different configurations to occur is larger than the hadron traversal time. The important message to be drawn is that a sufficiently energetic composite particle interacts with a target as a beam of almost incoherent quark–gluon configurations, but not as the naively expected entity – the composite particle at rest, since the center of mass motion cannot be factorized.

This physics can be accounted for formally within the framework of the Feinberg and Pomeranchuk [2] and Good and Walker [3] method of description of diffractive processes in terms of scattering states. It is convenient to decompose the wave function of the energetic projectile h over the eigenstates of the scattering matrix S

$$|h>= \sum_n c_n|n>, \quad (2)$$

where

$$S|n>= d_n|n> . \quad (3)$$

The scattering eigenstates are not eigenstates of the Hamiltonian and in order to establish the nature of the hadron matter they correspond to one has to decompose the states $|n>$ over the eigenstates of QCD Hamiltonian. We shall postpone to a future paper the discussion of this interesting question. The scattering states $|n>$ acquire a phase as a result of the interaction with the target and therefore

$$S|h>= \sum d_n c_n|n> . \quad (4)$$

Inelastic processes will occur if the phases d_n of the different states $|n>$ are different. By selecting certain final states, it is possible to enhance the role of some quark–gluon configurations in a projectile. This is the essence of the idea [4] on how to search for the color transparency phenomenon in quasielastic reactions (see Ref. [5] for a review and related references).

It is widely known by now that, at least within the eikonal approximation, configurations of different transverse spatial size r are scattering eigenstates and are characterized by different cross sections

$$\sigma \approx \frac{r^2}{< r^2 >} < \sigma > \tag{5}$$

for $r^2 \lll r^2 >$, as a result of color screening. Here $< \sigma >$ is the measured mean value of cross section. For a review and earlier references on the theoretical and experimental aspects in support of Eq. (5) in perturbative and nonperturbative QCD, and of the emerging physical picture, see review [5]. This feature of the interaction can be used to disentangle the contributions arising from different spatial size configurations. The coherence of the configurations which satisfy Rel. (1), but having different interactions with the target, can be formally accounted for by introducing a distribution over the values of the cross section $P(\sigma)$ [7], instead of an average value of the cross section only.

The number of frozen configurations $|n>$ during a collision depends on the initial energy, see Rel. (1), and this number increases with the energy of the projectile. In the case of a nucleon projectile this physics becomes relevant at $E_N > 40$ GeV, for $A \approx 200$ in the nucleus rest frame (at CERN and RHIC energies), which follows from using $M_n = m_{N^*}$ in Rel. (1) – the first excitation of a nucleon [6]. The states with smaller invariant masses are not important, since the pion is a pseudogoldstone particle of the spontaneously broken chiral symmetry in QCD.

The analysis of the experimental data on the diffractive hadron production in pp–scattering [7] and on the inelastic shadowing correction to the total cross section of pd–scattering [8] allows the extraction of the second moment (dispersion) of this cross section probability distribution $P(\sigma)$

$$w = \frac{< \sigma^2 >}{< \sigma >^2} - 1 \approx 0.25 \dots 0.5. \tag{6}$$

Different numbers in Rel. (6) correspond to different energies of the projectile. Such large values are in agreement with the quark models of the nucleon [6]. The analysis of the diffractive hadron production off a deuteron [8] leads to

$$\frac{<\sigma^3>}{<\sigma>^3} - 1 \approx 3w \tag{7}$$

By definition

$$<\sigma^n>= \frac{\int P(\sigma)\sigma^n d\sigma}{\int P(\sigma)d\sigma}. \tag{8}$$

These numbers were extracted from the theoretical analysis of phenomena in the several hundred GeV energy range of a proton projectile. There are strong indications that the fluctuations of cross sections increase with energy at least up to the RHIC energies [8]. Data clearly demonstrate that the σ–distribution is rather wide. We want to use these large fluctuations, already observed experimentally, to analyse the possibility of producing a beam of excited baryon rich matter. The idea is that the multiple scattering processes are determined by moments of the cross sections and as a consequence the interacting nucleon has effectively a larger than average spatial size. In particular, the cross–section fluctuations lead to an enhancement of the fluctuations of the number of NN–subcollisions and hence to larger fluctuations in transverse energy E_t [8] in comparison with those computed in an independent NN–collisions description with a constant (nonfluctuating) cross section [9]. Using the dispersion of the nucleon–nucleon cross sections measured at FNAL and ISR in single diffractive processes, it has been found [8] that this effect contributes significantly to the broadening of the E_t tail found by NA34 at CERN [10]. This leads us to the conclusion that the conventional use of large E_t as a trigger for centrality in heavy ion collisions, selects at the same time in the wave function of interacting nucleon a quark–gluon configurations of larger than average spatial size. We shall call such a configuration a huskion. On the other hand, the absence of the spectator nucleons in the event as the trigger for centrality of the collision, puts no restrictions on the wave function of the interacting nucleons.

Let us estimate the magnitude of the expected effects. For the sake of simplicity of presentation we shall consider only the properties of the projectile. However, a similar picture is valid for the target nuclei as well. In an independent particle model of the nucleus, the probability of a nucleon to be in a larger than average size configuration during a NN–collision, i.e. to be a huskion, is given by the formula:

$$p_1(\sigma_0) = \frac{\int_{\sigma_0}^{\infty} \sigma P(\sigma) d\sigma}{<\sigma>}. \quad (9)$$

σ_0 characterizes the size/huskiness of an interacting nucleon. In the case of a trigger for n NN–collisions the probability of formation of a huskion is

$$p_n(\sigma_0) = \frac{\int_{\sigma_0}^{\infty} P(\sigma)\sigma^n d\sigma}{<\sigma^n>}. \quad (10)$$

Examples of such triggers are $E_t \gg < E_t >$, $< E_t^n >$, etc. Rel. (10) for $n = 2$ is equivalent to Eq. (12) of ref. [8], obtained by using the Abramovsky, Gribov and Kancheli rules [11] for the calculation of inclusive spectrum. We shall use the cross section distributions extracted in Ref. [8] for the case when the dispersion of the cross section in Rel. (6) is minimal and therefore we shall underestimate the role of fluctuations. A simple numerical calculation shows that for $n \geq 3$

$$p = (p_n - p_1) \geq 0.2 \dots 0.3 \quad (11)$$

for $\sigma_0 > < \sigma >$, see Fig. 1. Thus at given impact parameter at least one nucleon is large with the probability $pN \approx 1$, where N is the number of nucleons at the same impact parameter (for heavy nuclei $N \approx A^{1/3} \approx 6$). In order to get a rough idea of the consequences of cross section fluctuations, let us consider an oversimplified model, where the probability for a nucleon to be a huskion is p and N nucleons at the same impact parameter are distributed uniformly. The probability P_k, to have at least k huskions at a given impact parameter is given by the following formula

$$P_k = \sum_{r=k}^{N} \binom{N}{r} p^r (1-p)^{N-r}, \quad (12)$$

$$P_1 = 1 - (1-p)^N, \quad P_2 = 1 - (1-p)^N - Np(1-p)^{N-1}, \dots.$$

Since the selection of the central collisions requires at least one nucleon at a given parameter to be large, the probability to have k huskions ($k > 2$) is given by the conditional probability P_k/P_1. It is easy to check that $P_2/P_1 \approx 0.5$ for $p \approx 0.2 \dots 0.3$ and $N = 6$. Obviously, this oversimplified model underestimates the number of huskions at the same impact parameter, since it neglects the correlations between 3 and more particles (which are large, see Rel. (7)), the larger radius of a huskion, etc.

From the presented arguments, it follows that the effective mass of a huskion depends on energy. At CERN nuclear beam energies according to Rel. (1) this effective mass is in between m_N and $\approx m_{N^*}$. Consequently, the physical object under consideration is a nucleus, where instantaneously many nucleon centers are excited. The question arises whether they form a network. This is the standard problem of the percolation theory, see review [12]. (We want to stress that the dynamics of internucleon interactions within the nucleus is not considered here – only collective effects due to the selection of central collisions). To account for the geometry of percolation one can decompose the instantaneous configuration over the normal nucleons and huskions and model nuclei as a system of spheres, which may be huskions with probability p and normal nucleons with probability $1 - p$. Results quoted in Ref. [12] show that an infinite cluster arises whenever $p > 0.2 - 0.3$. Our nontrivial observation is that percolation exists also for relatively large values of σ_0 also, see Fig. 1. The bottom line is that central nucleus–nucleus collisions provide a natural method to select clusters of superdense nuclear matter in nuclei.

In analogy with condensed matter physics, it is reasonable to expect that a system with randomly distributed excited centers corresponds to a new phase, if the probability of percolation between huskions is significant. This is a phase transition where the order parameter is p. Near the critical point, when $p \leq p_c$ and neglecting finite number nucleon effects, one can show that the formation probability of a network has a $(p - p_c)^a$–type of singularity [12], where $a < 0$ and p_c is the critical percolation probability. Also, the correlation length between two huskions increases with the distance d between them as $(p - p_c)^b$, where $b < 0$ as well. When percolation occurs, a fractal spatial distribution of the huskions appears. By applying the Abramovsky, Gribov and Kancheli rules [11] one can come to the conclusion that the final state hadron distributions near the percolation phase transition should reveal the fractal properties of the beam. It will be interesting to apply also an intermittency analysis [13] of the hadron distributions, in the central heavy ion collisions at CERN, RHIC and LHC energies, in order to search for such a fractal.

The percolation characterizes the coherence of the quark–gluon orbits in different nucleons, i.e. correlations between nucleons at different impact parameters. We want to emphasize that in normal nuclei such coherent effects are suppressed by a large barrier factor. In the high energy processes discussed above, this barrier factor is absent. However, we do not claim that either the projectile or target nuclei spontaneously undergo a phase transition into a superdense state. The answer to this question requires a theoretical analysis of the overlap integrals between the scattering and hadron states.

One of the methods of investigation of such nuclear beams would be the comparison of the parton distributions of the beam in the case of central collisions with large E_t as a trigger and with the trigger for lack of nucleons–spectators and/or with the parton distributions in a nucleon, measured in pp-collisions with a high transverse energy E_t trigger. This amounts to the determination of the ratio of the structure functions of a nucleus and a nucleon defined as

$$R_i(x, Q^2, b) = \frac{D_i^{beam}(x, Q^2, b)}{AD_i^N(x, Q^2)}.$$

Here x is the usual Bjorken variable, b is the impact parameter in nucleus–nucleus collisions and i = valence quark, sea quark or gluon. In the following we shall restrict our analysis to central nucleus–nucleus collisions only, i.e. $b \approx 0$.

The shadowing at small x should be more pronounced than in the case of the structure function of ordinary nuclei, due to the larger radius of huskions, i.e. $R_i < 1$ for $2xm_N r_{NN} < 1$. Here r_{NN} is the mean internucleon distance in nuclei. In this case, two qualitative effects follow from the analysis of Ref. [14]. The point where the depletion of parton distributions is changed by an enhancement at larger x, $x \approx 2m_N r_{NN}$, should move to the right with increasing energy. Besides, the absolute value of the enhancement of the parton distributions R_i should also increase with energy. This follows directly from the approach of Ref. [14], which is based on the calculation of nuclear shadowing at small x and the reconstruction of parton distributions of a nucleus at larger x, using the exact sum rules for the baryon charge and the total momentum of a nucleon. Thus, the prediction is that the gluon and valence quark distributions at $x \approx 2m_N r_{NN}$

should increase with energy in central heavy ion collisions, barring the well understood Q^2 evolution phenomenon. This enhancement of the gluon and valence quark distributions can be checked in high p_T phenomena, in the bottom quark production etc.

Other possible effects depend on the origin of the percolation phase transition. Let us consider one option, which is in line with the data on deep inelastic lepton scattering off nuclei. Experimental data found an enhancement of the valence quark distribution at $x \approx 2m_N r_{NN}$, when compared to that of a free nucleon and no enhancement for the sea quark distribution in nuclei [15, 16]. This enhancement of valence quark– and gluon–, but not of antiquark–distributions in nuclei has been interpreted in [14] as a manifestation of the color screening phenomenon and as a tendency towards a phase transition into a quark–gluon state, when increasing the nucleon density. If so even more peculiar behaviour is expected for the valence quark distributions in the beam at $x > 0.3$. In Ref. [17] it was shown that the probability of small spatial size quark–gluon configurations in a bound nucleon decreases with nuclear density as a consequence of the color screening phenomenon. Consequently, if the number of huskions in the beam is comparable with the atomic number of the beam, then $R_v(x > 0.3) \leq 1$ and it should decrease with energy, due to the color screening effects and the well understood evolution with Q^2. However, at sufficiently large energies and large n the overlap integral with the nucleon system is small (such a situation corresponds to the dissolution of the nucleons in the dense matter) one may expect the opposite effect, i.e. an increase of the distribution with density, once the well understood evolution of distributions with Q^2 has been taken into account. Correspondingly, at such energies R_v should begin to increase with energy. For the valence quarks, with larger transverse intrinsic momenta, this effect should be even more pronounced, since this selection enhances the role of the small size configurations in the interacting nucleon.

A more detailed investigation of the huskion matter may be feasible at RIHC and LHC energies, where the fragmentation regions of the colliding nuclei are well separated. In particular, if the nuclear beam in central heavy ion collisions corresponds to the formation of a quark–gluon phase, one may expect the occurrence of the effects discussed in connection with the conventional quark–gluon plasma (see Ref. [18] for a review and additional references): the enhancement of strange and charmed particle production for $x < A^{-1/3}$, etc. We want to emphasize that the theoretical tools used in this paper give no clue on either the origin of the percolation phase transition or its uniqueness. However, in the opinion of the present authors, the existence of such a phase transition looks like an inevitable logical conclusion, in the light of emerging current physical picture of high–energy collisions.

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REFERENCES

1. H.D.Politzer, Phys. Rev. Lett. **30**, 1346 (1973); D.J.Gross and F.Wilczek, Phys. Rev. Lett. **30**, 1343 (1973).

2. E.L.Feinberg and I.Y.Pomeranchuk, Suppl.Nuovo Cimento **III**, 652(1956).

3. M.L.Good and W.D.Walker, Phys.Rev. **120**, 1857 (1960).

4. A.H.Mueller, in "Proceedings of the seventh Rencontre de Moriond", ed. J.Tran Than Van, Edtions Frontieres, Gif–sur–Yvette, France, (1982), p. 13; S.J.Brodsky, in "Proceedings of the thirteenth symposium on multiparticle dynamics", eds. W.Kittel, W.Metzger and A.Stergman, World Scientific, Singapore (1982), p. 963.

5. L.L.Frankfurt and M.Strikman, Prog.Part.Nucl.Phys. **27**, 135 (1991).

6. L.L.Frankfurt and M.Strikman, Phys.Rev.Lett. **66**,2289 (1991).

7. H.Miettinen and J.Pumplin, Phys.Rev. **D18**, 1696 (1978); Phys.Rev.Lett. **42**, 204 (1979).

8. H.Heiselberg, G.Baym, B.Blattel, L.Frankfurt and M.Strikman, Phys.Rev.Lett. **76**, Nov. (1991) ; B.Blattel, G.Baym, L.Frankfurt, H.Heiselberg and M.Strikman,in *Proc. Quark Matter 91*, Nucl.Phys. A, (in press) and preprint, University of Illinois at Urbana-Champaign, P92–6–66 (1992).

9. G.Baym, G.Friedman, H.Heiselberg, in *Proc. Brookhaven HIPASS Workshop*, March 1990, ed. O.Hansen.

10. J.Schukraft et al.(NA34), Nucl.Phys **A498**, 67c (1989); T.Akesson et al; (Helios Collaboration), Nucl.Phys. **B353**, 1(1991).

11. V.Abramovski, V.Gribov and O.Kancheli, Sov.J.Nucl.Phys. **18**, 308 (1973).

12. D.Stauffer,Phys.Rep.**54**,1 (1979).

13. A.Bialas and R.Peschanski, Nucl.Phys. **B273** 703 (1986), *ibid* **B308**, 857 (1988); A.Bialas and R.C.Hwa, Phys.Lett. **B 253**, 436 (1991).

14. L.Frankfurt, M.Strikman and S.Liuti, Phys. Rev. Lett. **65**, 1725 (1991).

15. J.J.Aubert et al, Phys. Lett. **B123**, 275 (1983); R.Arnold et al, Phys.Rev. Lett. **52**, 1431 (1984).

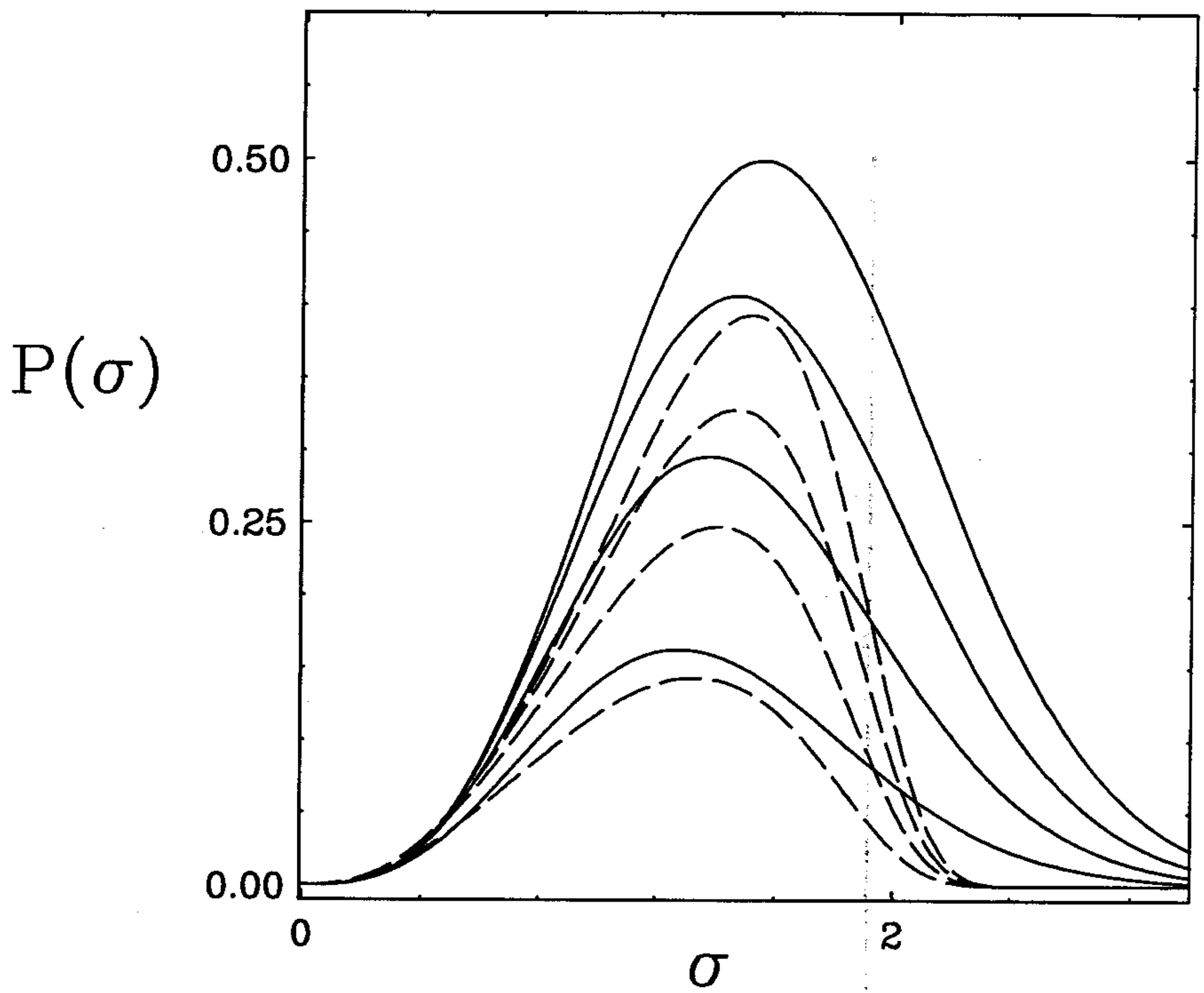
16. D.Alde et al, Phys. Rev. Lett. **64**, 2479 (1990).

17. L.Frankfurt and M.Strikman, Nucl. Phys. **B250**, 146 (1985); L.Frankfurt and M.Strikman, Phys. Rep. **160**, 235 (1988).

18. H.Satz, in *Heavy Ion Physics at Very High Energies*, CERN, Geneva (1991).

FIGURE CAPTION

Fig.1 The σ_0–dependence of the probability $p = p_n(\sigma_0) - p_1(\sigma_0)$, see Eq. (11), for $n = 2, \ldots, 5$, $(p_{n+1}(\sigma_0) > p_n(\sigma_0))$. Two parametrizations of the cross–section probability distribution $P(\sigma) = N\sigma/(\sigma + a\sigma_1)\exp[-(\sigma/\sigma_1 - 1)^m/\Omega^m]$, see Ref. [8], with $m = 2$, $a = 1$, $\sigma_1 = 0.63$, $\Omega = 1.5$ (solid lines) and $m = 10$, $a = 1$, $\sigma_1 = 0.16$, $\Omega = 11$ (dashed lines) have been used. Both these parametrizations are characterized by a dispersion $w = 0.25$, see Rel. (6). σ is in units of $< \sigma >$.


P(σ)
0.50
0.25
0.00
0
2
σ