HADRONIZATION OF SUPERCOOLED QGP IN ULTRA-RELATIVISTIC HEAVY ION COLLISIONS

V. MAGAS[1,2], L.P. CSERNAI[3,4], M.I. GORENSTEIN[1,5], L.L. JENKOVSZKY[1]

[1]**Bogolyubov Institute for Theoretical Physics, Nat. Acad. Sci. of Ukraine**
(14b, Metrolohichna Str., Kyiv 03143, Ukraine),

[2]**Physics Departament, Instituto Superior Tecnico**
(Ave. Rovisco Pais, 1049-001 Lisbon, Portugal),

[3]**Department of Physics, University of Bergen**
(Allegaten 55, N-5007, Norway),

[4]**Research Institute for Particle and Nuclear Physics**
(P.O. Box 49, 1525 Budapest, Hungary),

[5]**University of Frankfurt**
(Robert-Mayer Str., 8-10, D-60054 Frankfurt am Main, Germany)

UDC 539

The important role of two-particle correlation measurements in modern ultra-relativistic heavy ion collisions is discussed as well as the latest results from RHIC and SPS. Possible hadronization of supercooled QGP created in such reactions is studied within the Bjorken hydrodynamic model. Such a hadronization should be a very fast shock-like process, which, hadronization coincides or is shortly followed by freeze-out, could explain a part of the HBT puzzle, i.e., the flash-like particle emission. HBT data also show that the total expansion time before freeze-out is very short ($\sim 6-10$ fm/c). Here, we discuss the question of supercooled QGP and the timescales of the reaction.

1. HBT Radii and Their Interpretation

Two-particle interferometry has become a powerful tool for studying the size and duration of particle production from elementary collisions (e^+e^-, pp and $p\bar{p}$) to heavy ions like Au+Au at RHIC or Pb+Pb at SPS [1—3]. For the case of nuclear collisions, the interest mainly focuses on the possible transient formation of a deconfined state of matter. This could affect the size of the region from where the mesons (mostly pions) are emitted as well as the time for particle production.

The two-pion correlation function measures the coincidence probability $P(\mathbf{p}_1, \mathbf{p}_2)$ of two (identical) bosons with momenta $\mathbf{p}_1$, $\mathbf{p}_2$ relative to the probability of detecting uncorrelated particles from different events,

$$C_2(\mathbf{p}_1, \mathbf{p}_2) = \frac{P(\mathbf{p}_1, \mathbf{p}_2)}{P(\mathbf{p}_1)P(\mathbf{p}_2)}\,. \tag{1}$$

We denote $K^\mu = (p_1^\mu + p_2^\mu)/2$ and $q^\mu = p_1^\mu - p_2^\mu$. Then assuming that the particle source is chaotic and sufficiently large,

$$C_2(\mathbf{p}_1, \mathbf{p}_2) = 1 + \frac{\left|\frac{1}{(2\pi)^3}\int_{\Sigma_{\rm FO}} d\Sigma K\ \exp\left[i\,\Sigma q\right]\, f\left(\frac{uK}{T}\right)\right|^2}{E_1\,\frac{dN}{d^3\mathbf{p}_1}\ E_2\,\frac{dN}{d^3\mathbf{p}_2}}\,, \tag{2}$$

where [20]

$$E\,\frac{dN}{d^3\mathbf{p}} = \frac{1}{(2\pi)^3}\int\limits_{\Sigma_{\rm FO}} d\Sigma p\ f\left(\frac{up}{T}\right) \tag{3}$$

is the single inclusive momentum distribution, $f(x) = (e^x-1)^{-1}$ is the pion distribution, and u^μ is the fluid 4-velocity. The integrals run over some unknown *ad hoc* freeze-out (FO) hypersurface — therefore, the further results are model dependent.

The normalized correlation functions are fitted by a Gaussian:

$$C_2(\vec{q}) = 1 + \lambda\exp(-R^2_{\rm long}q^2_{\rm long} - $$
$$-R^2_{\rm side}q^2_{\rm side} - R^2_{\rm out}q^2_{\rm out} - 2R^2_{\rm outlong}q_{\rm out}q_{\rm long}), \tag{4}$$

with $R_{\rm long}$, $R_{\rm side}$, $R_{\rm out}$ being the Gaussian source radii, or HBT radii, and λ the correlation strength. The cross-term $R^2_{\rm outlong}$ appears as a consequence of space-time correlations in non-boost-invariant systems.

In a very simple case of the Bjorken longitudinal boost-invariant model, the measured HBT radii can receive a straightforward physical meaning. For the Bjorken cylinder geometry, it is useful to restrict the consideration to particles emitted at midrapidity, $K^z = q^z = 0$. Rotational symmetry around the z-axis in central collisions makes it possible to choose the avera-

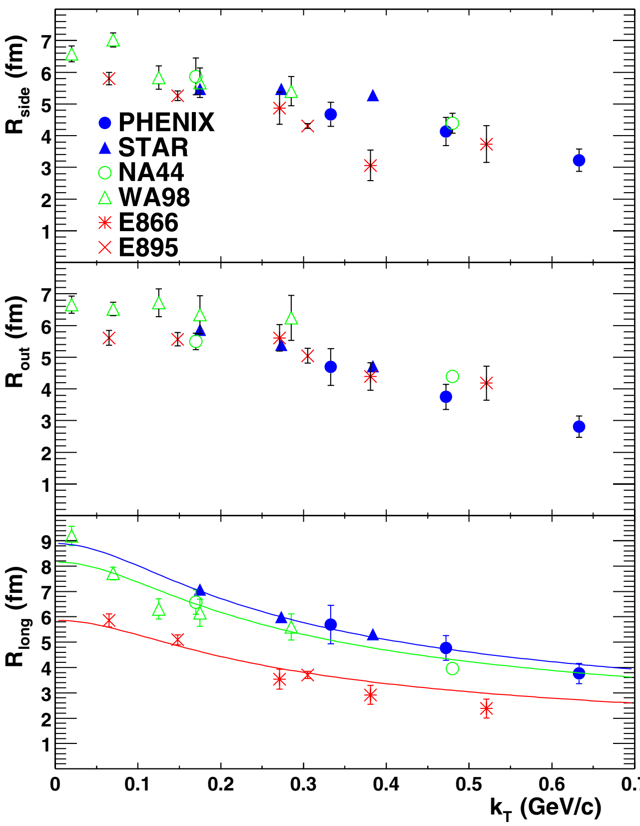

Fig. 1. HBT radii for pion pairs as a function of k_T measured at mid-rapidity for various energies from E895 ($\sqrt{s_{NN}}$ = 4.1 GeV), E866 ($\sqrt{s_{NN}}$=4.9 GeV), NA44, WA98 ($\sqrt{s_{NN}}$=17.3 GeV), STAR, and PHENIX ($\sqrt{s_{NN}}$=130 GeV). The bottom plot includes fits to $A/\sqrt{m_T}$ for each energy region. The data are for π^- results except for the NA44 results, which are for π^+. From [PHENIX Collaboration, Phys. Rev. Lett. **88**, 192302 (2002)]

ge transverse momentum as $\mathbf{K}_\perp = (K, 0, 0)$, and consequently, $C_2(K, q_{\rm out}, q_{\rm side})$ is a function of three independent variables only. The so-called out and side projections of the relative momenta are $\mathbf{q}_{\rm out} = (q_{\rm out}, 0, 0)$, $\mathbf{q}_{\rm side} = (0, q_{\rm side}, 0)$. As shown in [4, 3] the width, $1/R_{\rm side}$, of the correlation function in $q_{\rm side}$ is a measure of the transverse decoupling or the freeze-out radius, while the width $1/R_{\rm out}$ of the $q_{\rm out}$ correlation function is also sensitive to the duration of hadronization, $\Delta\tau$:

$$R^2_{\rm out} \approx R^2_{\rm side} + v^2 \Delta\tau^2 \quad . \tag{5}$$

In the general case for a longitudinal boost-invariance system, $R_{\rm long}$ can be connected to the lifetime τ_s of the system, the time elapsed between the onset of expansion and the kinetic freeze-out, $\tau_{\rm FO}$, by [6]

$$R_{\rm long}(m_t) = \tau_{\rm FO}(T_{\rm FO}/m_t)^{\frac{1}{2}}, \tag{6}$$

where $T_{\rm FO}$ is the freeze-out temperature, m_t is the transverse mass.

In the general case (if we do not assume boost-invariance), the physical interpretation of the HBT radii — $R_{\rm long}$, $R_{\rm side}$, $R_{\rm out}$ — is much more less certain. Nevertheless, this is a believe in the HBT physics society that the above-mentioned interpretation within theBjorken model captures the main tendency and, therefore, $R_{\rm out}$ is used to be associated with the geometric freeze-out radius. From transverse mass spectrum of the $R_{\rm long}$, one extracts $\tau_{\rm FO}$, and $R_{\rm out}/R_{\rm side}$ is used as a measure of the hadronization time.

Calculations show that a strong first-order QCD phase transition within continuous hydrodynamical expansion would lead to long lived gradually hadronizing QGP. Such a system behaviour would manifest itself as a large $R_{\rm out}/R_{\rm side}$ ratio. This scenario is not supported by experimental data [5], which show that $R_{\rm out}/R_{\rm side}$ is close to 1 or even smaller than 1(!)[1] – see Fig. 1. This suggest very fast flash-like hadronization.

Comparing the recent data [5] from RHIC with SPS data, one finds a "puzzle" [7]: all the HBT radii are pretty similar although the center of mass energy is changed by an order of magnitude (see Fig. 1). Discussions at "Quark Matter 2002" [8] lead to the conclusion that the duration of particle emission, as well as the lifetime of the system before freeze-out, appears to be shorter than the predictions of most of the models at the physics market.

An alternative possibility, discussed in [12—16], is the hadronization from the supercooled QGP (sQGP). This is expected to be a very fast shock-like process. If the hadronization from sQGP coincides with freeze-out, like it was assumed in [15], then this could explain a part of the HBT puzzle, i.e. the flash-like particle emission ($R_{\rm out}/R_{\rm side} \approx 1$). In this work, we are asking the following question — can the hadronization from sQGP explain also the another part of the HBT puzzle, i.e. a very short ($\sim$ 6[17] – 10[8, 18] fm/s) expansion time before freeze-out?

2. Shock Hadronization of the sQGP

Relativistic shock phenomena were widely discussed with respect to their connection to high-energy heavy ion collisions (see, for example, [19]). In thermal equilibrium by admitting the existence of the sQGP

[1]These results principally question our understanding of the space-time evolution of $A + A$!

and the superheated hadronic matter (HM) we have essentially richer picture of discontinuity-like transitions than in standard compression and rarefaction shocks. The system evolution in relativistic hydrodynamics is governed by the energy-momentum tensor $T^{\mu\nu} = (\epsilon + p)u^\mu u^\nu - pg^{\mu\nu}$ and conserved charge currents (in our applications to heavy ion collisions, we consider only the baryonic current nu^μ). They consist of local thermodynamical fluid quantities (the energy density ϵ, pressure p, baryonic density n) and the collective four-velocity $u^\mu = \sqrt{1-\mathbf{v}^2}(1,\mathbf{v})$. Continuous flows are the solutions of the hydrodynamical equations:

$$\partial_\mu T^{\mu\nu} = 0 \ , \quad \partial_\mu n u^\mu = 0 \ , \tag{7}$$

with specified initial and boundary conditions. These equations are nothing more than the differential form of the energy-momentum and baryonic number conservation laws. Along with these continuous flows, the conservation laws can also be realized in the form of discontinuous hydrodynamical flows which are called shock waves and satisfy the following equations:

$$T_o^{\mu\nu} d\sigma_\nu = T^{\mu\nu} d\sigma_\nu, \quad n_0 u_0^\mu d\sigma_\mu = n u^\mu d\sigma_\mu, \tag{8}$$

where $d\sigma^\mu$ is the unit 4-vector normal to the discontinuity hypersurface. In Eq. (8), the zero index corresponds to the initial state ahead of the shock front and quantities without index are the final state values behind it. A general derivation of the shock equations (valid for both space-like and time-like normal vectors $d\sigma^\mu$) was given in [21].

The important constraint on transitions (8) (thermodynamical stability condition) is the requirement of nondecreasing entropy (s is the entropy density):

$$s u^\mu d\sigma_\mu \geq s_0 u_0^\mu d\sigma_\mu \ . \tag{9}$$

To simplify our consideration and make our arguments more transparent, we consider only one-dimensional hydrodynamical motion. To study the shock transitions at the surface with space-like (s.l.) normal vector (we call them s.l. shocks), one can always choose the Lorentz frame where the shock front is at rest. Then $d\sigma^\mu = (0,1)$ at the surface of shock discontinuity, and Eq. (8) in this (standard) case becomes:

$$T_0^{01} = T^{01} \ , \ T_0^{11} = T^{11} \ , \ n_0 u_0^1 = n u^1 \ . \tag{10}$$

Solving Eq. (10), one obtains [14]:

$$v_0^2 = \frac{(p-p_0)(\epsilon+p_0)}{(\epsilon-\epsilon_0)(\epsilon_0+p)} \ ,$$

$$v^2 = \frac{(p-p_0)(\epsilon_0+p)}{(\epsilon-\epsilon_0)(\epsilon+p_0)} \ , \tag{11}$$

and the well known Taub adiabat (TA) [22]

$$n^2X^2 - n_0^2X_0^2 - (p-p_0)(X+X_0) = 0 \ , \tag{12}$$

where $X \equiv (\epsilon+p)/n^2$.[2]

For discontinuities on a hypersurface with a time-like (t.l.) normal vector $d\sigma^\mu$ (we call them t.l. shocks), one can always choose another convenient Lorentz frame ("simultaneous system") where $d\sigma^\mu = (1,0)$. Equation (8) is then

$$T_0^{00} = T^{00} \ , \ T_0^{10} = T^{10} \ , \ n_0 u_0^0 = n u^0. \tag{13}$$

Solving Eq. (13), we find

$$\tilde{v}_0^2 = \frac{(\epsilon-\epsilon_0)(\epsilon_0+p)}{(p-p_0)(\epsilon+p_0)}, \quad \tilde{v}^2 = \frac{(\epsilon-\epsilon_0)(\epsilon+p_0)}{(p-p_0)(\epsilon_0+p)} \ , \tag{14}$$

where we use the "$\sim$" sign to distinguish the t.l. shock case (14) from the standard s.l. shocks of (11). Another relation contains only the thermodynamical variables. It appears to be identical to the TA of Eq. (12). Eqs. (14) and (11) are connected to each other by simple relations [14]:

$$\tilde{v}_0^2 = \frac{1}{v_0^2} \ , \ \ \tilde{v}^2 = \frac{1}{v^2} \ . \tag{15}$$

These relations show that only one kind of transition can be realized for a given initial state and final state. The physical regions $[0,1)$ for v_0^2, v^2 (11) and for $\tilde{v}_0^2, \tilde{v}^2$ (14) can be easily found in the (ϵ–p)-plane [14]. For a given initial state (ϵ_0, p_0), they are shown in Fig. 2. For supercooled initial QGP states, the TA no longer passes through the point (ϵ_0, p_0) and new possibilities of t.l. shock hadronization transitions to regions III and VI in Fig. 2 appear.

[2] It has been shown in a series of works [23], that freeze-out through the space-like hypersurface leads to a nonequilibrium post FO distribution.

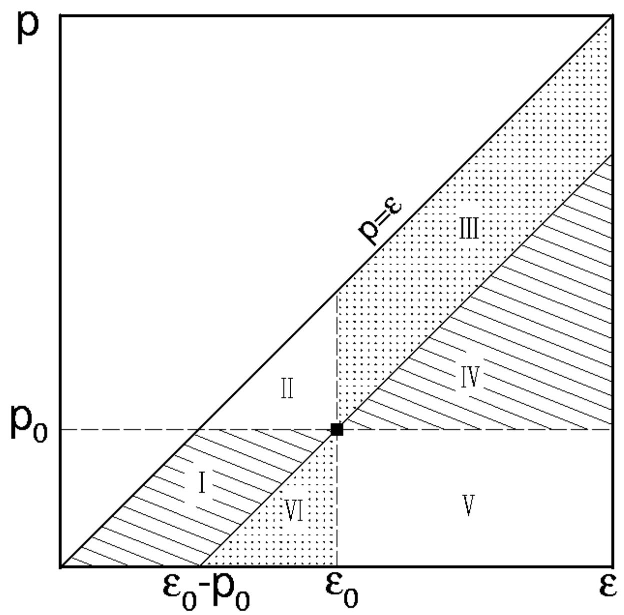

Fig. 2. Possible final states in the (energy density–pressure)-plane for shock transitions from the initial state (ϵ_0, p_0). I and IV are the physical regions for s.l. shocks, III and VI for t.l. shocks. II and V are unphysical regions for both types of shocks. Note that only states with $p \leq \epsilon$ are possible for any physical Equation of State in the relativistic theory

3. Hadronization of the sQGP within Bjorken Hydrodynamics

For a study of the expanding QGP, we have chosen a framework of the one-dimensional Bjorken model [24] (actually our principal results will not change if we use a $3D$ Bjorken model). Within the Bjorken model, all the thermodynamical quantities are constant along constant proper time curves, $\tau = \sqrt{t^2 - z^2}$ =const. The important result of Bjorken hydrodynamics (which assumes a perfect fluid) is that the evolution of the entropy density, is independent of the Equation of State (EoS), namely

$$s(\tau) = \frac{s(\tau_{\text{init}})\tau_{\text{init}}}{\tau} \, . \tag{16}$$

In the Bjorken model, the natural choice of a freeze-out hypersurface is the τ =const hypersurface, where the normal vector is parallel to the Bjorken flow velocity, $v = z/t$. Thus, $d\sigma^\nu = (1,0)$ in the rest frames of each fluid element. This leads to the simple solution of the t.l. shock equations (13):

$$\tilde{v}^2 = \tilde{v}_0^2 = 0 \; , \;\; \epsilon = \epsilon_0, \; n = n_0, \;\; p \neq p_0 \; . \tag{17}$$

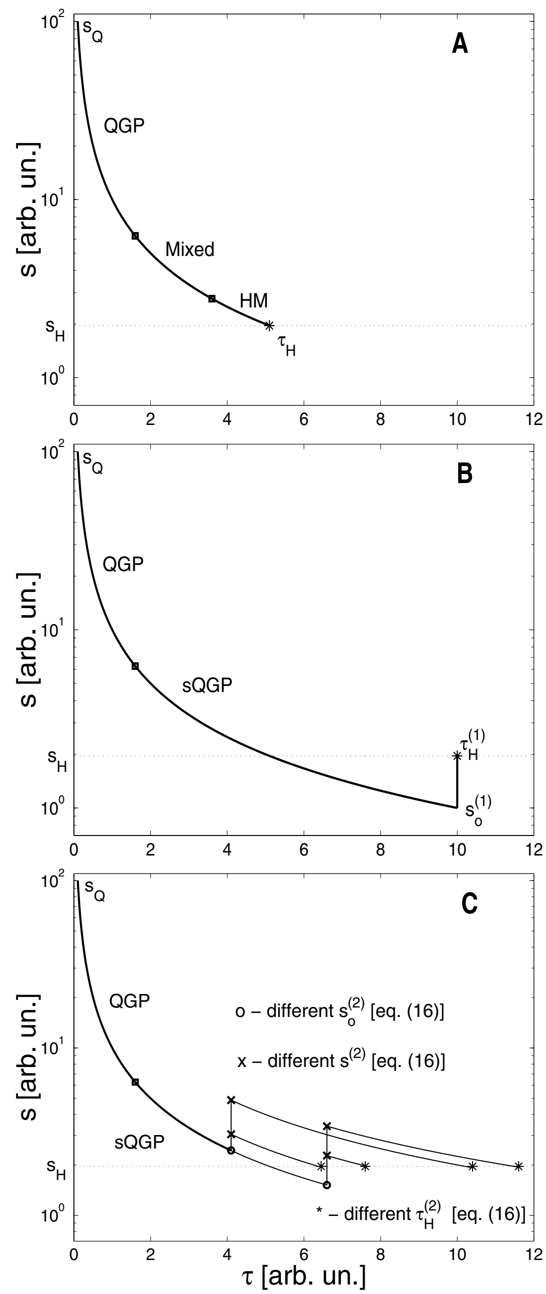

Fig. 3. Different ways for a system to go from Q state (s_Q) to H state (s_H) are presented on $\{s, \tau\}$ plane. Subplot A shows continuous expansion, which takes time τ_H, Eq. (19). Subplot B presents flash-like particle emission, i.e. simultaneous hadronization and freeze-out; which takes time $\tau_H^{(1)}$, Eq. (20). Subplot C shows several possibilities according to scenario 2 with shock-like hadronization into superheated HM. Time $\tau_H^{(2)}$, Eq. (22), can be smaller or larger than $\tau_H^{(1)}$, depending on details of the EoS, but always larger than τ_H

The entropy condition (9) is reduced to

$$s \geq s_0 \ . \tag{18}$$

Now let us try to answer the main question of this work — can QGP expansion with t.l. shock hadronization of supercooled state be faster than the hadronization through the mixed phase? The initial state is given at the proper time $\tau_{\rm init} \equiv \tau_Q$, when the local thermal equilibrium is achieved in the QGP state $Q \equiv (\epsilon_Q, p_Q, s_Q)$. The final equilibrium hadron state is also fixed, by experiment or otherwise, as $H \equiv (\epsilon_H, p_H, s_H)$. For the continuous expansion given by Eq. (16), the proper time for the $Q \to H$ transition is (see Fig. 3 — subplot A):

$$\tau_H = \frac{s_Q \tau_Q}{s_H} \ . \tag{19}$$

If our system enters the sQGP phase and the particle emission is flash-like, i.e. the system hadronizes and freezes out at the same time, then Eq. (16) is also valid all the time with final t.l. shock transition to the same H state. We call this as a scenario number one (see Fig. 3 — subplot A). Our system should go into the supercooled phase to the point where $\epsilon_0^{(1)} = \epsilon_H, n_0^{(1)} = n_H$, as it is required by Eq. (17). At this point, our sQGP has entropy density $s_0^{(1)}$. Its value depends on the EoS, but the t.l. shock transition is only possible if $s_0^{(1)} \leq s_H$ according to Eq. (18). Thus, for the proper time of the $Q \to H$ transition according to the first scenario, we have:

$$\tau_H^{(1)} = \frac{s_Q \ \tau_Q}{s_0^{(1)}} \geq \tau_H \ . \tag{20}$$

We can also study a scenario number two when our system supercools to the state $(\epsilon_0^{(2)}, p_0^{(2)}, s_0^{(2)})$, then hadronizes to a superheated HM state $(\epsilon^{(2)}, p^{(2)}, s^{(2)})$, and then this HM state expands to the same freeze-out state $H \equiv (\epsilon_H, p_H, s_H)$. (see Fig. 3 — subplot C). At the point of the shock transition, one has:

$$\tau_0^{(2)} = \frac{s_Q \ \tau_Q}{s_0^{(2)}} \ , \tag{21}$$

Then we have a t.l. shock transition satisfying Eq. (17), and, following the HM branch of the hydrodynamical expansion, we find:

$$\tau_H^{(2)} = \frac{s^{(2)} \ \tau_0^{(2)}}{s_H} = \frac{s_Q \tau_Q}{s_H} \frac{s^{(2)}}{s_0^{(2)}} \geq \tau_H \ , \tag{22}$$

since $s^{(2)} \geq s_0^{(2)}$ due to the non-decreasing entropy condition (18). In this second scenario, the value of the entropy density $s_0^{(2)}$ of sQGP can be both smaller and larger than the HM final value s_H. Depending on details of the EoS, the proper time $\tau_H^{(2)}$ (22) of the $Q \to H$ transition can also be smaller as well as larger than $\tau_H^{(1)}$ (20), but always larger than τ_H.

Conclusions

The conclusion of our analysis seems to be a rather general one: the system's evolution through the supercooled phase and time-like shock hadronization can not be shorter than a continuous expansion within the perfect fluid hydrodynamics independently of the details of EoS and the parameter values of the initial, Q, and final, H, states. Although we may achieve a flash-like particle emission in such a way, supported by the HBT data, the expansion time becomes longer, making it harder to reproduce the experimental HBT radii.

And, in fact, we do not see any physical process which would help us to achieve a shorter freeze-out time than the minimal one coming from the (very fast) Bjorken expansion via thermal and phase equilibrium. Any delay in the phase equilibration (see assignment 9 in [25]) or/and any dissipative process in our system lead to the entropy production, what increases the time needed to reduce the entropy density to $s_0 \leq s_H$.

The construction of a full reaction model, which simultaneously describes data on two-particle interferometry, hadron spectra, and hadron abundances is a formidable task which is still ahead of us.

We acknowledge the financial support from the Humboldt Foundation and INTAS grant 00-00366.

1. *Shuryak E.*// Phys. Lett. **B44** (1973) 387; *Gyulassy M., Kauffmann S. K., WilsonL. W.*// Phys. Rev. **C20** (1979) 2267; *Sinyukov Yu.M.*// Nucl. Phys. **A498** (1989) 151, **A556** (1994) 589; *Heinz U. W., Jacak B. V.*// Ann. Rev. Nucl. Part. Sci. **49** (1999) 529; *Weiner R. M.*// Phys. Rept. **327** (2000) 249.

2. *Pratt S.*// Phys. Rev. Lett. **53** (1984) 1219; Phys. Rev. **D33** (1986) 1314.

3. *Bertsch G., Gong M., Tohyama M.*// Phys. Rev. **C37** (1988) 1896; *Bertsch G.*// Nucl. Phys. **A498** (1989) 173.

4. *Pratt S.*// Phys. Rev. **C49** (1994) 2722,; Phys. Rev. **D33** (1986) 1314.

5. *RHIC data*/ C. Adler et al. [STAR Collaboration]; Phys. Rev. Lett. **87** (2001) 082301; *Adcox et al.* [PHENIX Collaboration]; **88** (2002) 192302; *SPS data*/ S. V. Afanasev et al. [NA49 Collaboration]; Nucl. Phys. **A698** (2002) 104; *Appelshäuser H.*// hep-ph/0204159.
6. *Makhlin A.N., Sinyukov Yu.M.*// Z. Phys. **C39** (1988) 69.
7. *Gyulassy M.* nucl-th/0106072.
8. *Pratt S.*// Talk at the Quark Matter 2002, July, Nantes, France.
9. *Rischke D., Gyulassy M.*// Nucl. Phys. **A608** (1996) 479.
10. *Teaney D., Lauret J., Shuryak E.V.*// Phys. Rev. Lett. **86** (2001) 4783; nucl-th/0110037.
11. *Csernai L.P., Kapusta J.I.*// Phys. Rev. Lett. **69** (1992) 737; Phys. Rev. **D46** (1992) 1379.
12. *Csörgő T., Csernai L.P.*// Phys. Lett. **B333** (1994) 494.
13. *Csernai L.P., Mishustin I.N.*// Phys. Rev. Lett. **74** (1995) 5005.
14. *Gorenstein M.I., Miller H.G., Quick R.M., Ritchie R.A.*// Phys. Lett. **B340** (1994) 109.
15. *Keranen A., Mannien J., Csernai L.P., Magas V.*// nucl-th/0205019.
16. *Jenkovszky L.L., Kampfer B., Sysoev V.M.*// Z. Phys. **C48** (1990) 147.
17. *Csörgő T., Ster A.*// nucl-th/0207016.
18. *Panitkin S.*// Proc. XXXII Intern. Sumposium on Multiparticle Dynamics, Sept. 7—13, 2002, Alushta, Ukraine.
19. *Stock R.*// Phys. Repts. **135** (1986) 259; *Stöcker H., Greiner W.*// Ibid. **137** (1986) 277; *Bugaev K.A., Gorenstein M.I., Kämpfer B., Zhdanov V.I.*// Phys. Rev. **D40** (1989) 2903.
20. *Cooper F., Frye G., Schonberg E.*// Phys. Rev. **D11**, (1975) 192.
21. *Csernai L.*// Zh. Eksp. Teor. Fiz. **92** (1987) 379; Sov. Phys. JETP **65** (1987) 216.
22. *Taub A.H.*// Phys. Rev. **74** (1948) 328.
23. *Anderlik Cs. et al*// Phys. Rev. **C59** (1999) 388; Ibid. (1999) 3309; Phys. Lett. **B459** (1999) 33; *Magas V.K. et al.*// Heavy Ion Phys. **9** (1999) 193; Nucl. Phys. **A661** (1999) 596.
24. *Bjorken J.D.*// Phys. Rev. **D27** (1983) 140.
25. *Csernai L.P.* Introduction to Relativistic Heavy Ion Physics.— New York: Wiley, 1994.

АДРОНІЗАЦІЯ ПЕРЕОХОЛОДЖЕНОЇ КВАРК-ГЛЮОННОЇ ПЛАЗМИ В УЛЬТРАРЕЛЯТИВІСТСЬКИХ ЗІТКНЕННЯХ ВАЖКИХ ІОНІВ

В. Магас, Л.П. Чернаі, М.Й. Горенштейн, Л.Л. Єнковський

Р е з ю м е

Обговорено важливу роль вимірювань двочастинкових кореляцій в ультрарелятивістських зіткненнях важких іонів та останні дані, отримані на найсучасніших прискорювачах RHIC та SPS. Досліджено можливу адронізацію переохолодженої кварк-глюонної плазми, створеної в таких реакціях, у бйоркенівській моделі. Така адронізація була б надзвичайно швидким шокоподібним процесом, який, якщо заморожування збігається в часі або настає відразу за адронізацією, може пояснити частину “загадки HBT”, а саме спалахоподібне випромінювання частинок. Дані HBT також показують, що загальний час розширення надзвичайно короткий. У цій роботі ми обговорюємо питання, пов"язані з адронізацією переохолодженої кварк-глюонної плазми та характерними часами реакції.

АДРОНИЗАЦИЯ ПЕРЕОХЛАЖДЕННОЙ КВАРК-ГЛЮОННОЙ ПЛАЗМЫ В УЛЬТРАРЕЛЯТИВИСТСКИХ СТОЛКНОВЕНИЯХ ТЯЖЕЛЫХ ИОНОВ

В. Магас, Л.П. Чернаи, М.И. Горенштейн, Л.Л. Енковский

Р е з ю м е

Обсуждается важная роль измерений двухчастичных корреляций в ультрарелятивистских столкновениях тяжелых ионов и последние результаты, полученные на самых современных ускорителях RHIC и SPS . Изучается возможная адронизация переохлажденной кварк-глюонной плазмы, образовавшейся в таких реакциях, в бьеркеновской модели. Такая адронизация была бы очень быстрым шокоподобным процессом, который, если замораживание совпадает или следует вскоре после адронизации, может объяснить часть “загадки HBT”, а именно вспышкообразное испускание частиц. HBT-данные также указывают на очень короткое общее время расширения. В данной работе мы обсуждаем вопросы, связанные с адронизацией переохлажденной кварк-глюонной плазмы и характерными временами реакции.