Possible Frictionless Regime for
Ultra-High Temperature Amorphous Matter

Zotin K.-H. Chu

1Centre for Distribution, P.O. Box 39, Toudiban, Road Xihong, Urumqi 830000, China

Abstract

The almost frictionless transport of the very-high temperature amorphous matter which resembles the color glass condensate (possibly having much of their origin in the RHIC heavy ion collisions) in a confined annular tube with transversely corrugations is investigated by using the verified transition-rate model and boundary perturbation method. We found that for certain activation volume and energy there exist possible frictionless states which might be associated with the perfect fluid formation during the early expansion stage in RHIC Au+Au collisions. We also address the possible similar scenario in LHC Pb+Pb collisions considering the possible perfect fluid formation in ultra-high temperature transport of amorphous matter.

KEYWORDS: Phenomenological Models, Statistical Methods, Quantum Dissipative Systems, Thermal Field Theory, Boundary perturbation.

1 Introduction

The success of ideal hydrodynamics [1-7], in explaining bulk of the data in Au+Au collisions at RHIC (cf., e.g., [8-9]), has led to a paradigm that in Au+Au collisions a nearly perfect fluid is created. There is no a priori reason why ideal hydrodynamics works so well at RHIC.

The relation between the inverse Reynolds number \( R^{-1} \) and the attenuation length \( L_s = (4/3)\eta/(e + p) \) becomes \( R^{-1} \approx L_s/\tau \) in the Bjorken flow case \( (\eta \) is the shear viscosity, \( e \) is the energy density, \( p \) is the pressure, \( \tau \approx \sim 0.6 \text{ fm/c} \) is the expansion rate) [10-11]. \( R^{-1} \) is found to be very small from the blast-wave fitting with the above correction term (cf. [10]). This suggests that the hadronic fluid in heavy ion collisions is nearly a perfect one.

The perfect fluid can also be understood from the form of the stress-energy tensor. In the approximation where the Lagrangian is \( \mathcal{L} = P(X) \), with \( P(X) \) being the kinetic function, the stress-energy tensor has the form

\[
T_{\mu\nu} \propto -P(X)g_{\mu\nu} + 2P'(X)u_\mu u_\nu,
\]

where \( u_\mu = \partial_\mu \phi \equiv u, \phi \) is a (real) scalar. This has the form of a stress-energy tensor for a perfect fluid with 4-velocity \( u_\mu \) [12]. The fact that \( u_\mu \) is a gradient means that the flow is irrotational: \( \nabla \times u = 0 \). Note that the defining property of a perfect fluid (fluid with zero viscosity [13]) is that at each point there is a frame in which the stress-energy tensor has the form \( T_{\mu\nu} = \text{diag}(-\rho, p, p, p) \) and in ideal hydrodynamics, the mean free path of particles is assumed to be zero.
The interesting (relevant to our borrowed approach here) issue is the possible existence of the color glass condensate (CGC) (cf. e.g., [13,14,7,8]). Au+Au collisions are then collisions of two sheets of colored glass, with the produced quarks and gluons materializing at a time given by the inverse of the saturation momentum. The proposed color glass condensate was thought to be a possible precursor state to the quark-gluon plasma (QGP). To be specific, at very high energies, multi-particle production in QCD is generated by small $x$ partons in the nuclear wavefunctions ($x$ is the fraction of the longitudinal momentum carried by the parton). These partons have properties best described as a color glass condensate (CGC) [5]. When two sheets of CGC collide in a high-energy heavy-ion collision, these partons are released and create energy densities an order of magnitude above the energy density required for the crossover from hadronic to partonic degrees of freedom. This matter, at early times after a heavy-ion collision, is a coherent classical field, which expands, decays into nearly on shell partons and may eventually thermalize to form a quark-gluon plasma. Because it is formed by melting the frozen CGC degrees of freedom and is the non-equilibrium matter preceding the QGP, this matter is called the glasma [5,13] (a non-equilibrium gluonic state between the collision moment and equilibrated QGP [6]).

As remarked in [13], the color glass condensate is the matter associated with the wee partons of a high energy hadronic wavefunction. This matter has very high energy density. It has properties similar to Bose-Einstein condensates (there is a characteristic momentum scale, termed the saturation scale, below which the gluon density saturates. This effect sets in when $x$ becomes small and the associated gluon wave length increases to nuclear dimensions. In such a regime gluons may interact and form a coherent state reminiscent of a Bose-Einstein condensate) and to spin glasses. It was motivated based on observations in deep inelastic scattering. The idea is that as one goes to higher and higher energies, there are more and more gluons in the hadron wavefunction. They have to squeeze together, and highly occupy phase space, making a highly coherent high density system of gluons. This matter controls the high energy limit of hadronic scattering and provides the initial conditions for the matter made in such collisions. Thus the almost perfect fluid phenomena (or frictionless transport) observed during the early stage (expansion) for Au+Au collisions in RHIC could be dominated by small $x$ (gluons) (so that the geometrical scaling occurs) [13,14]. Note that the gluon population at low $x$ is not an incoherent superposition of nucleon structure functions but is limited with increasing $Au$ by nonlinear gluon-gluon fusion resulting from the overlap of gluons from several nucleons in the plane of the nucleus transverse to the collision axis.

In fact due to the nearly perfect fluid (very low viscosity or flow resistance being almost zero) characteristics associated to a possible new state of matter created in Au+Au collisions at RHIC some researchers argued that this property (the (kinematic) shear viscosity of this nearly perfect fluid has been determined and found to be at least a factor of 4 smaller than that of the superfluid $^4$He [2]) could be called high temperature superfluidity because the fluid of quarks in the super-high temperature range flows better than the superfluid $^4$He in the extremely low
temperature region (cf. [2,4]). With above facts we can have in mind the possible link between the (color) glassy condensate (with the role of gluons) and their almost frictionless transport (of the fluid or liquid). This latter situation is similar to the proposed superglass reported quite recently [15,16].

Meanwhile we know that at low temperatures gases condense into the liquid (or solid) state. In a liquid transport is no longer governed by the motion of individual molecules or composite (condensed) particles. As noticed in [1], Eyring proposed that momentum transport of condensed composite particles or molecules (liquid-like) is due to processes that involve the motion of vacancies [17-19]. These processes can be viewed as thermally activated transitions in which a molecule or a cluster moves from one local energy minimum to another. Eyring’s approach involves the Planck constant $h$ and the appearance of $h$ is related to Eyring’s assumption that the collision time of the molecules is $h/(k_B T)$, the shortest timescale in a liquid ($k_B$ is the Boltzmann constant, $T$ is the temperature). One important observation is: Eyring’s approach has been successfully applied to the study of transport of glassy matter at a wide range of temperatures [17-19]. Possible frictionless transport of glassy matter under specific environment were also reported recently [16,20].

Above mentioned analogy or observed facts is the primary motivation for our present study. Our concern is associated with the possible superfluidity (or perfect fluid) formation after shear-thinning (i.e., the viscosity diminishes with increasing shear rate.). While superflow (or perfect fluid flow) in a state of matter possessing a shear modulus might initially seem untenable, experimental claims for precisely this phenomenon in solid $^4$He now abound (cf. [15,16,20].

In this paper, to demonstrate the nearly frictionless transport of CGC and identify the possible transition (or critical) temperature, we shall adopt the verified transition-rate-state model [17-19] to study the transport of CGC (presumed to be amorphous) within a corrugated annular tube (shell-like). The possible nearly frictionless states due to strong shear-thinning will be relevant to the perfect fluid formation at the early stage for Au+Au collisions in RHIC as mentioned above. To obtain the law of shear-thinning matter for explaining the too rapid annealing at the earliest time, because the relaxation at the beginning was steeper than could be explained by the bimolecular law, a hyperbolic sine law between the shear (strain) rate $\dot{\gamma}$ and shear stress $\tau$ has been proposed and the close agreement with experimental data was obtained [17-19]. This model has sound physical foundation from the thermal activation process [1,17-19] (a kind of quantum tunneling which relates to the matter rearranging by surmounting a potential energy barrier was discussed therein). With this model we can associate the (shear-thinning) fluid with the momentum transfer between neighboring clusters on the microscopic scale and reveals the possible microscopic interaction in the relaxation of flow with dissipation (the momentum transfer depends on the activation (shear) volume $V^* \equiv V_h$ which is associated with the center distance between atoms and is equal to $k_B T/\tau_0$ ($\tau_0$ a constant with the dimension of stress).

To consider the more realistic but complicated confined boundary conditions in the interfaces of
the annular tube (shell-like), however, we will use the boundary perturbation technique [16,20] to handle the presumed wavy-roughness along the interfaces of the confined annual tube. To obtain the analytical and approximate solutions, here, the roughness is only introduced in the radial or transverse direction. The relevant boundary conditions along the wavy-rough surfaces will be prescribed below. We shall describe our approach after this section: Introduction with the focus upon the transition-rate approach and boundary perturbation method. The approximate expression of the transport is then demonstrated at the end. Finally, we will illustrate our results into three figures and give discussions therein.

2 Theoretical Formulations

Researchers have been interested in the question of how (complex) matter responds to an external mechanical load. External loads cause transport, in Newtonian or various types of non-Newtonian ways. Amorphous matter, composed of polymers, metals, or ceramics, can deform under mechanical loads, and the nature of the response to loads often dictates the choice of matter in various applications. The nature of all of these responses depends on both the temperature and loading rate [1].

To the best knowledge of the author, the simplest model that makes a prediction for the rate and temperature dependence of shear yielding is the rate-state model of stress-biased thermal activation [17-19]. Structural rearrangement is associated with a single energy barrier $E$ that is lowered or raised linearly by an applied stress $\sigma$: $R_\pm = \nu_0 \exp[-E/(k_BT)] \exp[\pm \sigma V^*/(k_BT)]$, where $\nu_0$ is an attempt frequency and $V^*$ is a constant called the 'activation volume'. In amorphous matter, the transition rates are negligible at zero stress. Thus, at finite stress one needs to consider only the rate $R_+$ of transitions in the direction aided by stress. The linear dependence will always correctly describe small changes in the barrier height, since it is simply the first term in the Taylor expansion of the barrier height as a function of load. It is thus appropriate when the barrier height changes only slightly before the system escapes the local energy minimum. This situation occurs at higher temperatures; for example, Newtonian transport is obtained in the rate-state model in the limit where the system experiences only small changes in the barrier height before thermally escaping the energy minimum. As the temperature decreases, larger changes in the barrier height occur before the system escapes the energy minimum (giving rise to, for example, non-Newtonian transport). In this regime, the linear dependence is not necessarily appropriate, and can lead to inaccurate modeling. To be precise, at low shear rates ($\dot{\gamma} \leq \dot{\gamma}_c$), the system behaves as a power law shear-thinning material while, at high shear rates, the stress varies affinely with the shear rate. These two regimes correspond to two stable branches of stationary states, for which data obtained by imposing either $\sigma$ or $\dot{\gamma}$ exactly superpose.

We shall consider a steady transport of the amorphous matter (CGC) in a wavy-rough annu-
lar tube of $r_1$ (mean-averaged inner radius) with the inner interface being a fixed wavy-rough surface: $r = r_1 + \epsilon \sin(k\theta + \beta)$ and $r_2$ (mean-averaged outer radius) with the outer interface being a fixed wavy-rough surface: $r = r_2 + \epsilon \sin(k\theta)$, where $\epsilon$ is the amplitude of the (wavy) roughness, $\beta$ is the phase shift between two walls, and the roughness wave number: $k = 2\pi/L$ ($L$ is the wavelength of the surface modulation in transverse direction).

Firstly, this amorphous matter (composed of composite condensed particles, say, quarks and gluons) can be expressed as [16,18,20] $\ddot{\gamma} = \dot{\gamma}_0 \sinh(\tau/\tau_0)$, where $\dot{\gamma}$ is the shear rate, $\tau$ is the shear stress, $\tau_0 = 2k_B T/V_h$, and $\dot{\gamma}_0(\equiv C_k k_B T \exp(-\Delta E/k_B T)/h)$ is with the dimension of the shear rate; here $C_k \equiv 2V_h/V_m$ is a constant relating rate of strain to the jump frequency ($V_h = \lambda_2 \lambda_3 \lambda$, $V_m = \lambda_2 \lambda_3 \lambda_1$, $\lambda_2 \lambda_3$ is the cross-section of the transport unit on which the shear stress acts, $\lambda$ is the distance jumped on each relaxation, $\lambda_1$ is the perpendicular distance between two neighboring layers of particles sliding past each other), accounting for the interchain co-operation required, $h$ is the Planck constant, $\Delta E$ is the activation energy. In fact, the force balance gives the shear stress at a radius $r$ as $\tau = -(r \delta \dot{G})/2$ [16,20]. $\delta \dot{G}$ is the net effective forcing along the transport (or tube-axis: $z$-axis) direction (considering $dz$ element).

Introducing the forcing parameter $\Phi = -(r_2/2\tau_0)\delta \dot{G}$ then we have $\dot{\gamma} = \dot{\gamma}_0 \sinh(\Phi r/r_2)$. As $\dot{\gamma} = -du/dr$ ($u$ is the velocity of the transport in the longitudinal ($z$-)direction of the annular tube), after integration, we obtain

$$u = u_s + \frac{\dot{\gamma}_0 r_2}{\Phi} \left[ \cosh \Phi - \cosh \left( \frac{\Phi r}{r_2} \right) \right],$$

(2)

here, $u_s(\equiv u_{\text{slip}})$ is the velocity over the (inner or outer) surface of the annular (cosmic) string, which is determined by the boundary condition. We noticed that a general boundary condition for transport over a solid surface [16,20] was

$$\delta u = L_s^0 \dot{\gamma} (1 - \frac{\dot{\gamma}}{\dot{\gamma}_c})^{-1/2},$$

(3)

where $\delta u$ is the velocity jump over the solid surface, $L_s^0$ is a constant slip length, $\dot{\gamma}_c$ is the critical shear rate at which the slip length diverges. Note that the slip (velocity) boundary condition above (related to the slip length) is closely linked to the mean free path of the particles together with a geometry-dependent factor (in low temperature regime it is the quantum-mechanical scattering of Bogoliubov quasiparticles which is responsible for the loss of transverse momentum transfer to the confined interfaces [21]). The value of $\dot{\gamma}_c$ is a function of the corrugation of interfacial energy.

With the slip boundary condition [16,20], we can derive the velocity fields and transport rates along the wavy-rough annular (cosmic) string below using the verified boundary perturbation technique [16,22] and dimensionless analysis. We firstly select $L_s^0$ to be the characteristic length scale and set $r' = r/L_s^0$, $R_1 = r_1/L_s^0$, $R_2 = r_2/L_s^0$, $\epsilon' = \epsilon/L_s^0$. After this, for simplicity, we drop all the primes. It means, now, $r$, $R_1$, $R_2$ and $\epsilon$ become dimensionless ($\Phi$ and $\dot{\gamma}$ also follow). The wavy interfaces are prescribed as $r = R_2 + \epsilon \sin(k\theta)$ and $r = R_1 + \epsilon \sin(k\theta + \beta)$ and the presumed steady transport is along the $z$-direction (annulus-axis direction).
2.1 Boundary Perturbation

Along the outer interface (the same treatment below could also be applied to the inner interface), we have \( \dot{\gamma} = (du)/(dn) \) on interfaces. Here, \( n \) means the normal. Let \( u \) be expanded in \( \epsilon \):

\[
 u = u_0 + \epsilon u_1 + \epsilon^2 u_2 + \cdots,
\]

and on the boundary, we expand \( u(r_0 + \epsilon dr, \theta(= \theta_0)) \) into

\[
 u(r, \theta) = u(r_0, \theta) + \epsilon [du(r_0, \theta)] + \epsilon^2 [d^2u(r_0, \theta)] + \cdots = \{u_{\text{slip}} + \frac{\dot{\gamma} R_2}{\Phi} \text{[cosh} \Phi - \text{cosh}(\Phi_{R}/R_2)]\} \text{on interfaces, \quad } r_0 \equiv R_1, R_2;
\]

where

\[
 u_{\text{slip}} \text{on interfaces} = L_s^0 \{(1 - \frac{\dot{\gamma}^2}{\gamma_c})^{-1/2}\} \text{on interfaces,}
\]

Now, on the outer interface (cf., e.g., [16, 22])

\[
 \dot{\gamma} = \frac{du}{dn} = \nabla u \cdot \frac{\nabla (r - R_2 - \epsilon \sin(k\theta))}{\nabla (r - R_2 - \epsilon \sin(k\theta))} = [1 + \epsilon^2 \frac{k^2}{r^2} \cos^2(k\theta)] - \frac{1}{2} [u_r | (R_2 + \epsilon dr, \theta) -
\]

\[
 \frac{k}{r^2} \cos(k\theta) u_\theta | (R_2 + \epsilon dr, \theta)] = u_{0r} | R_2 + \epsilon [u_{1r} | R_2 + u_{0rr} | R_2 \sin(k\theta)] -
\]

\[
 \frac{k}{r^2} u_{0r} | R_2 \cos(k\theta)] + \epsilon^2 [- \frac{k^2}{2 r^2} \cos^2(k\theta) u_{0r} | R_2 + u_{2r} | R_2 + u_{1rr} | R_2 \sin(k\theta)] +
\]

\[
 \frac{1}{2} u_{0rr} | R_2 \sin^2(k\theta) - \frac{k}{r^2} \cos(k\theta) (u_{1\theta} | R_2 + u_{0\theta} | R_2 \sin(k\theta))] + O(\epsilon^3).
\]

Considering \( L_s^0 \sim R_1, R_2 \gg \epsilon \) case, we also presume \( \sin \Phi \ll \frac{\dot{\gamma}}{\gamma_0} \). With equations (2) and (6), using the definition of \( \dot{\gamma} \), we can derive the velocity field \( (u) \) up to the second order:

\[
 u(r, \theta) = -(R_2 \dot{\gamma}_0/\Phi) \{\cosh(\Phi_{R}/R_2) - \cosh \Phi [1 + \epsilon^2 \Phi^2 \sin^2(k\theta)/(2R_2^2)] +
\]

\[
 \epsilon \Phi \sinh \Phi \sin(k\theta)/R_2 \} + u_{\text{slip}} | r = R_2 + \epsilon \sin(k\theta).
\]

The key point is to firstly obtain the slip velocity along the boundaries or surfaces. After lengthy mathematical manipulations, we obtain the velocity fields (up to the second order) and then we can integrate them with respect to the cross-section to get the transport (volume flow) rate \( Q \), also up to the second order here:

\[
 Q = \int_0^{\theta_\Phi} \int_{R_1 + \epsilon \sin(k\theta + \beta)}^{R_2 + \epsilon \sin(k\theta)} u(r, \theta) r dr d\theta = Q_0 + \epsilon Q_{p0} + \epsilon^2 Q_{p2}.
\]

In fact, the approximate (up to the second order) net transport (volume flow) rate reads:

\[
 Q = \pi \dot{\gamma}_0 \{L_s^0 (R_2^2 - R_1^2) \sinh \Phi (1 - \frac{\sinh \Phi \Phi}{\gamma_c/\gamma_0})^{-1/2} + \frac{R_2}{\Phi} [(R_2^2 - R_1^2) \cosh \Phi - \frac{2}{\Phi} (R_2^2 \sinh \Phi -
\]

\[
 R_1^2 \cosh \Phi)]\}.
\]
$R_1 R_2 \sinh(\Phi R_2 / R_1) + 2 R_2^2 \Phi \tau (\cosh \Phi - \cosh(\Phi R_2 / R_1))] + \epsilon^2 \left( \frac{\pi}{2} u_{\text{slip}} (R_2^2 - R_1^2) + \right.

L^0 s^2 \tau_0 \sinh \Phi (1 + \frac{\sinh \Phi}{\gamma_c / \tau_0}) (-k^2 + \Phi^2)[1 - (R_1 / R_2)^2] + \frac{\pi}{2} \gamma_0 [R_1 \sinh(\Phi R_2 / R_1) - R_2 \sinh \Phi] -

\frac{\pi \gamma_0}{2} \frac{R_2}{\Phi} [\cosh \Phi - \cosh(\Phi R_2 / R_1)] + \frac{\pi}{4} \gamma_0 \Phi \cosh \Phi [R_2 - R_1^2 / R_2] +

\pi \gamma_0 \{[\sinh \Phi + L^0 s \cosh \Phi (1 + \frac{\sinh \Phi}{\gamma_c / \tau_0})][(R_2 - R_1 \cos \beta)] + \frac{\pi}{2} \gamma_0 \frac{R_2}{\Phi} \cosh \Phi +

L^0 s^2 \Phi^2 \frac{\cosh \Phi}{\gamma_c / \gamma_0} [1 - (R_1 / R_2)^2 \cosh \Phi]. \tag{7}

Here,

$u_{\text{slip}} = L^0 s \gamma_0 [\sinh \Phi (1 - \frac{\sinh \Phi}{\gamma_c / \tau_0})^{-1/2}]. \tag{8}$

### 3 Results and Discussions

We firstly check the roughness effect (or combination of curvature and confinement effects) upon the transport via strongly shearing because there are no available experimental data and numerical simulations for the same geometric configuration (annular tube with wavy corrugations in transverse direction). With a series of forcings : $\Phi \equiv -R_2 (\delta G)/(2 \tau_0)$, we can determine the enhanced shear rates ($d\gamma/dt$) due to gravity forcings. From equation (6), we have (up to the first order)

$$d\gamma \over dt = \frac{d\gamma_0}{dt} [\sinh \Phi + \epsilon \sin(k\theta) \frac{\Phi}{R_2} \cosh \Phi]. \tag{9}$$

The parameters are fixed below (the orientation effect : $\sin(k\theta)$ is fixed here). $r_2$ (the mean outer radius) is selected as the same as the slip length $L^0 s$. The amplitude of wavy roughness can be tuned easily. The effect of wavy-roughness is significant once the forcing ($\Phi$) is rather large (the maximum is of the order of magnitude of $\epsilon [\Phi \tanh(\Phi)/R_2]$).

If we select a (fixed) temperature, then from the expression of $\tau_0$, we can obtain the shear stress $\tau$ corresponding to above gravity forcings ($\Phi$) :

$$\tau = \tau_0 \sinh^{-1} [\sinh(\Phi) + \epsilon \sin(k\theta) \frac{\Phi}{R_2} \cosh(\Phi)]. \tag{10}$$

There is no doubt that the orientation effect ($\theta$) is also present for the amorphous matter. For illustration below, we only consider the maximum case : $|\sin(k\theta)| = 1$. We shall demonstrate our transport results below. The wave number of roughness in transverse direction is fixed to be 10 (presumed to be the same for both interfaces of the annular tube) here.

As the primary interest of present study is related to the possible frictionless transport or formation of superfluidity (presumed to be relevant to the CGC as mentioned in Introduction) due to strong shearing, we shall present our main results in the following. Note that, based on the
absolute-reaction-rate Eyring model (of stress-biased thermal activation), structural rearrangement is associated with a single energy barrier (height) $\Delta E$ that is lowered or raised linearly by a (shear) yield stress $\tau$. If the transition rate is proportional to the plastic (shear) strain rate (with a constant ratio: $C_0; \dot{\gamma} = C_0 R_t$, $R_t$ is the transition rate in the direction aided by stress), we have

$$\tau = 2\left[\frac{\Delta E}{V_h} + \frac{k_B T}{V_h} \ln\left(\frac{\dot{\gamma}}{C_0 \nu_0}\right)\right] \quad \text{if} \quad \frac{V_h \tau}{k_B T} \gg 1$$

(11)

where $\nu_0$ is an attempt frequency or transition rate, $C_0 \nu_0 \sim \dot{\gamma}_0 \exp(\Delta E/k_B T)$, or

$$\tau = 2\frac{k_B T}{V_h} \frac{\dot{\gamma}}{C_0 \nu_0} \exp(\Delta E/k_B T) \quad \text{if} \quad \frac{V_h \tau}{k_B T} \ll 1.$$ 

(12)

It is possible that the frictional resistance (or shear stress) can be almost zero (existence of $\tau \sim 0$) from above equations (say, equation (11) considering a sudden jump of the resistance). The nonlinear character only manifests itself when the magnitude of the applied stress times the activation volume becomes comparable or greater in magnitude than the thermal vibrational energy.

Normally, the value of $V_h$ is associated with a typical volume required for a microscopic shear rearrangement. Thus, the nonzero transport rate (of the condensed composite (quarks and gluons) system) as forcing is absent could also be related to a barrier-overcoming or tunneling for shear-thinning matter along the wavy-roughness (geometric valley and peak served as potential surfaces) in annular tubes when the wavy-roughness is present. Once the geometry-tuned potentials (energy) overcome this barrier, then the tunneling (spontaneous transport) inside wavy-rough annular tubes occurs. Now, we start to examine the temperature effect. We fix the forcing $\Phi$ to be 1 as its effect is of the order $O(1)$ for the shear rate. As the gravity forcing ($\delta G$) might depend on the temperature ($[\delta G] = 2\tau_0 \Phi/R_2$, $\tau_0 \equiv \tau_0(T)$, $V^*(\equiv V_h)$ is presumed to be temperature independent here for simplicity). Note that, according to [7], $V^* = 3V \delta \gamma/2$ for certain matter during an activation event [7], where $V$ is the deformation volume, $\delta \gamma$ is the increment of shear strain.

As the primary interest of present study is related to the possible phase transition [14-16] or formation of superfluidity (presumed to be relevant to the formation of dark matter mentioned in Introduction) due to strong shearing, we shall present our main results in the following. We performed intensive calculations or manipulations of related physical and geometric parameters, considering a hot big-bang universe [14] and examine what happens as it expands and cools through the transition temperature $T_c$. The selected temperature range and the activation energy follows this reasoning. Note that in unified models of weak and electromagnetic interactions $T_c$ is of the order of the square root of the Fermi coupling constant [14], $G_F^{1/2}$, i.e. a few hundred GeV. Thus the transition occurs when the universe is aged between $10^{-10}$ and $10^{-12}$ seconds and far above nuclear densities [23]. One possible high-temperature superfluidity formation (or nearly frictionless transport) regime [2,4] is demonstrated in Fig. 1. The activation energy
$(\Delta E)$ is $6 \times 10^{-10}$ Joule. In fact, all the results shown in this figure depend on $\dot{\gamma}_0$ and are thus very sensitive to $\Delta E$ (and $V_h$). Here $C_k = 2$ and the sudden jump of the shear stress (directly linked to the friction) occurring around $T \sim 10^{12}$ °K could be the transition temperature for the selected $\Delta E$ and $C_k$. There is a sudden friction drop around two orders of magnitude below $T \sim 2 \times 10^{12}$ °K ($V_h \sim 4 \times 10^{-11}$ m$^3$) and it is almost frictionless below $T \sim 10^{12}$ °K. The temperature regime we identify here is close to the critical (at the critical end point) as well as crossover (to the quark-gluon plasma (QGP) at zero baryon chemical potential,) temperature mentioned in [2] (cf. [1]).

The possible reasoning for this formation can be illustrated in Fig. 2. It could be due to the strong shearing driven by larger forcings along a confined tube. The shear-thinning (the viscosity diminishes with increasing shear rate) reduces the viscosity significantly. One possible outcome for almost vanishing viscosity is the nearly frictionless transport.

Based on the knowledge gained at RHIC, it should be possible to predict experimental results at the Large Hadron Collider (LHC), which will collide lead ions at much higher energy densities [24]. At this point it’s important to predict something about the LHC experiments [13,24] even the LHC will bring unanticipated discovery. Fig. 3 illustrates the possible (almost) perfect fluid formation or frictionless transport for ultra-high temperature transport of amorphous matter (possible color-glass-condensate [24]) considering Pb+Pb collisions at LHC. The activation energy is $10^{-9}$ Joule. We can observe there is a sudden friction (shear stress) drop around two orders of magnitude below $T \sim 7 \times 10^{17}$ °K ($V_h \sim 2.5 \times 10^{-6}$ m$^3$) and it is almost frictionless below $T \sim 3 \times 10^{17}$ °K.

4 Conclusions

To conclude in brief, we have obtained critical parameters for the possible perfect fluid formation or almost frictionless transport of ultra-high temperature amorphous matter (possible CGC or string-like composite (condensed) particles associated with voids) relevant to Au+Au collisions at RHIC as well as Pb+Pb collisions at LHC. These critical parameters depend strongly upon the temperature, activation energy and activation volume. We shall investigate other relevant issues [1,13,25,26] in the future.

References

[1] T. Schäfer and D. Teaney, Nearly perfect fluidity: from cold atomic gases to hot quark gluon plasmas, Rep. Prog. Phys. 72 (2009) 126001.

[2] R.A. Lacey, N.N. Ajitanand, J.M. Alexander, P. Chung, W.G. Holzmann, M. Issah, A. Taranenko, P. Danielewicz and H. Stöcker, Has the QCD critical point been signaled by observations at the BNL relativistic heavy ion collider? Phys. Rev. Lett. 98 (2007) 092301.
[3] L.P. Csernai, J.I. Kapusta, and L.D. McLerran, *Strongly interacting low-viscosity matter created in relativistic nuclear collisions*, Phys. Rev. Lett. 97 (2006) 152303.

[4] T. Csörgő, M.I. Nagy and M. Csanád, *New exact solutions of relativistic hydrodynamics*, J. Phys. G: Nucl. Part. Phys. 35 (2008) 104128.

[5] R. Venugopalan, *From plasma to quark-gluon plasma in heavy-ion collisions*, J. Phys. G: Nucl. Part. Phys. 35 (2008) 104003.

[6] E. Shuryak, *Quark-gluon plasma-new frontiers*, J. Phys. G: Nucl. Part. Phys. 35 (2008) 104044.

[7] I. Arsene, *et al.* (BRAHMS Collaboration), *Quark-gluon plasma and color glass condensate at RHIC? The perspective from the BRAHMS experiment*, Nuclear Physics A 757 (2005) 1.

[8] K. Adcox, *et al.* (PHENIX Collaboration), *Formation of dense partonic matter in relativistic nucleus-nucleus collisions at RHIC: Experimental evaluation by the PHENIX Collaboration*, Nucl. Phys. A 757 (2005) 184.

[9] M. Haack and A. Yarom, *Nonlinear viscous hydrodynamics in various dimensions using AdS/CFT*, JHEP 10 (2008) 063.

[10] T. Hirano, *Hydrodynamic models*, J. Phys. G: Nucl. Part. Phys. 30 (2004) S845.

[11] D. Teaney, *Effect of shear viscosity on spectra, elliptic flow, and Hanbury Brown-Twiss radii*, Phys. Rev. C 68 (2003) 034913.

[12] N. Arkani-Hamed, H.-C. Cheng, M.A. Luty, S. Mukohyamaae and T. Wiseman, *Dynamics of gravity in a Higgs phase*, JHEP 01 (2007) 036.

[13] L. McLerran, *From AGS-SPS and onwards to the LHC*, J. Phys. G: Nucl. Part. Phys. 35 (2008) 104001.

[14] D. Kharzeev, E. Levin and L. McLerran, *Parton saturation and N_part scaling of semi-hard processes in QCD*, Phys. Lett. B 561 (2003) 93.

[15] P. Phillips and A.V. Balatsky, *Cracking the superdolid*, Science 316 (2007) 1435.

[16] K.-H. W. Chu, *Nearly frictionless transport of glassy solid helium at low temperature regime*, Ann. Phys. (N.Y.) 323 (2008) 2474.

[17] R.E. Powell and H. Eyring, *Mechanism for the relaxation theory of viscosity*, Nature 154 (1944) 427.
[18] S. Glasstone, K.J. Laidler and H. Eyring, *The Theory of Rate Processes* (McGraw-Hill, New York, 1941).

[19] H. Eyring, D. Henderson, B.J. Stover and E.M. Eyring, *Statistical Mechanics and Dynamics* (Wiley, New York, 1964).

[20] Z.K.-H. Chu, *Possible rapid transport of glassy supersolid*. [Arxiv: 0707.2828].

[21] D. Einzel and J.M. Parpia, *Slip of quantum fluids*, *J. Low Temp. Phys.* **109** (1997) 1.

[22] W. K.-H. Chu, *Stokes slip flow between corrugated walls*, *ZAMP* **47** (1996) 591.

[23] T.W.B. Kibble, *Topology of cosmic domains and strings*, *J. Phys. A : Math. General* **9** (1976) 1387.

[24] M. Luzum and P. Romatschke, *Viscous hydrodynamic predictions for nuclear collisions at the LHC*, *Phys. Rev. Lett.* **103** (2009) 262302.

[25] G. Policastro, D.T. Son and A.O. Starinets, *The shear viscosity of strongly coupled N = 4 supersymmetric Yang-Mills plasma*, *Phys. Rev. Lett.* **87** (2001) 081601 [hep-th/0104066].

[26] L.H. Ford and N.F. Svaiter, *Quantum density fluctuations in classical liquids*, *Phys. Rev. Lett.* **102** (2009) 030602.
Figure 1. Comparison of calculated (shear) stresses using an activation energy $6 \times 10^{-10}$ J. There is a sharp decrease of shear stress around $T \sim 2 \times 10^{12}$ K. Below $10^{12}$ K, the transport of amorphous matter is nearly frictionless (cf. [1,2] for the critical temperatures related to Au+Au collisions at RHIC).
Figure 2. Increasing shear causes a local energy minimum to flatten until it disappears (energy barrier removal or quantum-like tunneling). The structural contribution to the shear stress is referred to shear-thinning.
Figure 3. Comparison of calculated (shear) stresses using an activation energy $10^{-9}$ J. There is a sharp decrease of shear stress around $T \sim 7 \times 10^{17}$ K. Below $3 \times 10^{17}$ K, the transport of amorphous matter is nearly frictionless (cf. [24] for the possible critical temperatures in Pb+Pb collisions at LHC).