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The effect of hadronization on the instabilities in an expanding parton plasma

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Abstract. The growth rate for instabilities in an expanding parton plasma is investigated by using a quasiparticle transport model including hadronization. The coupled Boltzmann equations for partons and pions with time dependent mean field masses and source terms are solved in the Bjorken boost invariant picture. Hadronization modifies the known instability in the parton plasma created by the mean field in two ways: In the beginning, hadronization increases the rate Γ of instability, but then $\Gamma \to 0$ when the hadronization is dominating the time evolution.

PACS. 25.75.-q Relativistic heavy-ion collisions – 12.38.Mh Quark-gluon plasma

Over the past several years the quasiparticle approximation has been found to be a good starting point for understanding relativistic heavy ion collisions and the related quark-gluon plasma (QGP) physics. For temperatures beyond the critical value T_c of the phase transition from a hadron gas to QGP, the equation of state obtained from lattice simulations [1] of quantum chromodynamics (QCD) can be reproduced by another system of noninteracting quasiparticles with an effective mass and a dynamical mean field potential which represents the remaining interactions which are not included in the effective mass [2–4]. Since the nonequilibrium effects play a crucial role in the evolution of a QGP produced in high energy nuclear collisions, the quasiparticle description of the lattice thermodynamics has been extended to investigating the parton plasma at nonequilibrium [5]. A transport model with a phase transition has recently been also discussed based on a Boltzmann equation at hadronic level [6]. Including the hadronization in the evolution of QGP is generally a non-trivial problem, yet it is an essential ingredient. In this letter we study how the hadronization affects the stability of an expanding parton plasma. We work in the framework of a quasiparticle transport model and treat hadronization in a phenomenological model. The aim is to get some qualitative insight into the role of hadronization.

Many authors have analysed the instability of nuclear matter [7–10] and quark plasma [11] in the linear response method. Recently the quasiparticle model [5], phenomenologically adjusted to the lattice data in equilibrium and therefore with correct deconfinement properties, has been used to study the spinodal and dynamical instabilities [12] of an expanding parton plasma. The spinodal instabilities related to a first-order phase transition are found to be rather slow, while the dynamical ones related to the rapid expansion are the dominant ones. These dynamical instabilities do not appear for the Nambu–Jona-Lasinio (NJL) model [13] even though it contains the mechanism of a first-order chiral phase transition. So the dynamical instabilities reflect confinement. Until now hadronization is not yet introduced into the discussion of instability. In the following we will first incorporate hadronization into transport model, and then analytically solve the transport equations in a longitudinally boost invariant expansion. We pay particular attention to the influence of the hadronization on the growth rate for instability.

To construct a quasiparticle transport model with coincides with the equation of state from lattice data in the thermodynamic limit, we assume that the quasiparticle masses and the mean field potential depend on temperature. Our quasiparticle system is assumed to consist of partons (with no distinction for quarks and gluons) and pions. The total pressure

$$P(m_p, m_\pi, T) = P_p(m_p, T) + P_\pi(m_\pi, T) - V(m_p, m_\pi) \ , \tag{1}$$

contains the partial pressures

$$P_p(m_p, T) = g_p \int \frac{d^3\mathbf{p}}{(2\pi)^3} \frac{\mathbf{p}^2}{3E_p} f_p(m_p, T, \mathbf{p}) \ ,$$
$$P_\pi(m_\pi, T) = g_\pi \int \frac{d^3\mathbf{p}}{(2\pi)^3} \frac{\mathbf{p}^2}{3E_\pi} f_\pi(m_\pi, T, \mathbf{p}) \ , \tag{2}$$

and the mean field potential V which is assumed to be a function of the quasiparticle masses m_p and m_π only. Effective QCD models like Friedberg-Lee [14] and NJL [13] in their quasiparticle limit lead to the same coupled system of (1) and (2), but give specific forms for V. In this letter we follow [5,12] and determine V from lattice data. In (2) T is the temperature of the plasma, g_p and g_π are degeneracies of partons and pions, and $E_p = \sqrt{m_p^2 + \mathbf{p}^2}$ and $E_\pi = \sqrt{m_\pi^2 + \mathbf{p}^2}$ are quasiparticle energies. Since the lattice simulation at finite baryon density is still poor, we do not consider the chemical potential dependence of the quasiparticle distributions. We neglect the difference between Bosons and Fermions by assuming distributions f_p and f_π in equilibrium to be Boltzmann functions

$$f^{eq}_{p/\pi}(E_{p/\pi}/T) = e^{-E_{p/\pi}/T} \; . \tag{3}$$

The masses m_p and m_π in the heat bath and the potential are not independent parameters in the model. They satisfy the so-called gap equations derived from the minimum property of the thermodynamic potential $\Omega = -P$,

$$\begin{aligned} \frac{\partial V}{\partial m_p}|_{m_\pi} + g_p \int \frac{d^3\mathbf{p}}{(2\pi)^3}\frac{m_p}{E_p} f_p(m_p, T) = 0 \; , \\ \frac{\partial V}{\partial m_\pi}|_{m_p} + g_\pi \int \frac{d^3\mathbf{p}}{(2\pi)^3}\frac{m_\pi}{E_\pi} f_\pi(m_\pi, T) = 0 \; . \end{aligned} \tag{4}$$

Using standard methods of thermodynamics one derives the total energy density of the system

$$\epsilon(m_p, m_\pi, T) = \epsilon_p(m_p, T) + \epsilon_\pi(m_\pi, T) + V(m_p, m_\pi) \tag{5}$$

with the quasi-particle energy densities given by

$$\begin{aligned} \epsilon_p(m_p, T) = g_p \int \frac{d^3\mathbf{p}}{(2\pi)^3} E_p f_p(m_p, T) \; , \\ \epsilon_\pi(m_\pi, T) = g_\pi \int \frac{d^3\mathbf{p}}{(2\pi)^3} E_\pi f_\pi(m_\pi, T) \; . \end{aligned} \tag{6}$$

The screening and dynamical masses for pions calculated on the lattice [15] are rather smooth over a large temperature region compared with the very steep change of the parton mass around the critical point of deconfinement phase transition. Therefore the temperature dependence, or in nonequilibrium the space-time dependence, of the pion mass will be left out. By substituting the constant pion mass $m_\pi = 140MeV$ into the pion energy density in (6), the dependence of the parton mass $m_p(T)$ and the dynamical potential $V(m_p(T))$ can be fixed from the requirement

$$\epsilon(m_p(T), m_\pi, T) = \epsilon_{lat}(T) \tag{7}$$

and the gap (4) for partons, where $\epsilon_{lat}(T)$ is the lattice energy density. The details of the determination of the parton mass and potential are similar to that in [5]. Corresponding to the lattice data [1] for 4 flavors, the noninteracting limit of QCD with 4 massless flavors yields the parton degeneracy $g_p = 62.8$.

After the equation of state of the quasiparticle model is fixed by knowing $m_p(T)$ we turn now to the extension of the model to the nonequilibrium case and describe the expanding parton plasma and the hadronization process. Under the assumption that for a given system the dynamics in nonequilibrium should be the same as that in equilibrium, one retains the gap (4) but replaces $m_p(T)$ and $f_p(\mathbf{p}, T)$ by $m_p(x)$ and $f_p(x, \mathbf{p})$ with $x = (\mathbf{x}, t)$. Then $m_p(x)$ can be calculated from $f_p(x, \mathbf{p})$ via (4). The distribution functions $f_{p/\pi}(x, \mathbf{p})$ for the density of partons/pions of momentum $\mathbf{p}$ at point $\mathbf{x}, t$ are governed by two Boltzmann kinetic equations,

$$\begin{aligned} \left(\partial_t + \frac{1}{E_p}\mathbf{p}\cdot\nabla - \frac{m_p}{E_p}\nabla m_p \cdot \nabla_p\right) f_p(x, \mathbf{p}) = -I_p \; , \\ \left(\partial_t + \frac{1}{E_\pi}\mathbf{p}\cdot\nabla - \frac{m_\pi}{E_\pi}\nabla m_\pi \cdot \nabla_p\right) f_\pi(x, \mathbf{p}) = I_\pi \; , \end{aligned} \tag{8}$$

where I_p and I_π on the right-hand sides are respectively loss and gain terms for the partons and pions associated with the hadronization process (we neglect thermalization processes). The above transport equations together with the parton gap equation determine simultaneously the space-time dependent distributions and the parton mass.

From the energy conservation in collisions, the parton and pion source terms obey the constraint

$$g_p \int \frac{d^3\mathbf{p}}{(2\pi)^3} E_p I_p - g_\pi \int \frac{d^3\mathbf{p}}{(2\pi)^3} E_\pi I_\pi = 0 \; . \tag{9}$$

The calculation of the hadronization rate within QCD is still not solved satisfactorily due to the non-perturbative nature. Up to now most studies of hadronization use phenomenological methods, for instance, the evaporation mechanism [16], Schwinger mechanism [17] of pair production, flux tube model [18], string fragmentation ([19]), cluster formation mechanisms ([20]) and effective Lagrangians [21]. As emphasised in the introduction, we focus our attention in this letter on the influence of the hadronization on the instabilities of the nonequilibrium parton plasma. Therefore we also employ a schematic model. We take here the relaxation time approximation and put all the hadronization properties into the two relaxation times θ_p and θ_π,

$$I_p = \frac{1}{\theta_p} f_p, \qquad I_\pi = \frac{1}{\theta_\pi} f_\pi \; . \tag{10}$$

In this way we can only discuss the transition from partons to pions, but not the back reaction. This is appropriate only for the expanding phase of the plasma. The characteristic scale θ for the hadronization is a few fm/c. The hadronization time θ_p for the transition of partons into mesons depends on the status of the system – there is no hadronization above the phase transition, and it becomes maximum at the phase transition. In the quasiparticle model, the state of the system reflects itself in the value for the effective mass. Therefore we make the following phenomenological ansatz for the relaxation time,

$$\frac{1}{\theta_p} = \frac{1}{\lambda}\frac{1}{2}\left(1+\tanh\frac{m_p - m_p^c}{\Delta m_p}\right) , \qquad (11)$$

where the space-time dependence of the hadronization is hidden in the mass dependence of the relaxation time. In our numerical calculation the time scale λ is taken as $1\ fm/c$, see [19]. The centre value $m_p^c = 270MeV$ is larger than the pion mass and thus satisfies the necessary mass condition for the hadronization. We choose $\Delta m_p = m_p^c/5$. While there are good reasons for the value of λ, values for m_p^c and Δm_p^c can only be termed reasonably. In the relaxation time approximation energy conservation (9) gives the pion relaxation time in terms of the parton relaxation time,

$$\theta_\pi = \frac{\epsilon_\pi}{\epsilon_p}\theta_p . \qquad (12)$$

The stability analysis of the coupled system of Boltzmann and gap equations proceeds as follows: We seek analytical solutions $f_p(x,\mathbf{p}), f_\pi(x,\mathbf{p})$ of these equations which describe an expanding system and analyse the response of the system to perturbations δf_p and δf_π. For the analytical solution we use Bjorken's boost invariant picture [22] which is valid in the central region of relativistic heavy ion collisions where a plateau structure may exist for the final rapidity distribution. Introducing the proper time $\tau = \sqrt{t^2 - z^2}$ and space-time rapidity $\eta = \frac{1}{2}\ln\frac{t+z}{t-z}$ instead of the time t and longitudinal coordinate z, and assuming that the effective parton mass and the quasiparticle densities in coordinate space (the integration of the phase-space distributions over momentum) depend on the proper time only, the transport (8) in the absence of the source terms have a scaling solution [23] $f(s, p_\perp)$ with $s = \frac{\tau}{\tau_0}m_\perp|\sinh(Y-\eta)|$, the initial proper time τ_0, the transverse mass $m_\perp(\tau) = \sqrt{m^2(\tau) + \mathbf{p}^2}$, and the particle rapidity $Y = \frac{1}{2}\ln\frac{E+p_z}{E-P_z}$. Transverse expansion is neglected. In order to fix the form of f we assume that at $\tau = \tau_0$ the system is in thermal equilibrium, i. e. given by (3). In the local rest frame which corresponds to the central slice of a relativistic heavy ion collision, $\eta = 0$ and $m_\perp \sinh Y = p_z$. Including the hadronization, the solutions of (8) in the local rest frame are just the above scaling solutions multiplied by a factor due to hadronization,

$$f_p(\tau,\mathbf{p}) = e^{-\int_{\tau_0}^{\tau}\frac{1}{\theta_p(\tau')}d\tau'} f_p^{eq}\left(\frac{\tilde{E}_p}{T}\right) ,$$
$$f_\pi(\tau,\mathbf{p}) = e^{\int_{\tau_0}^{\tau}\frac{1}{\theta_\pi(\tau')}d\tau'} f_\pi^{eq}\left(\frac{\tilde{E}_\pi}{T}\right) , \qquad (13)$$

with $\tilde{E}_{p/\pi} = \sqrt{(\frac{\tau}{\tau_0}p_z)^2 + p_\perp^2 + m_{p/\pi}^2(\tau_0)}$.

Substituting the scaling solution (13) into the gap equation for quarks and remembering that the dynamical potential V in nonequilibrium is the same as that extracted from the lattice data in equilibrium, we obtain in the local rest frame the time dependence of the parton mass shown in Fig.1a. In the numerical calculation the initial time and initial temperature are chosen as $\tau_0 = 1fm$ and $T = 270MeV$. As expected the quasiparticle mass increases with expansion time.

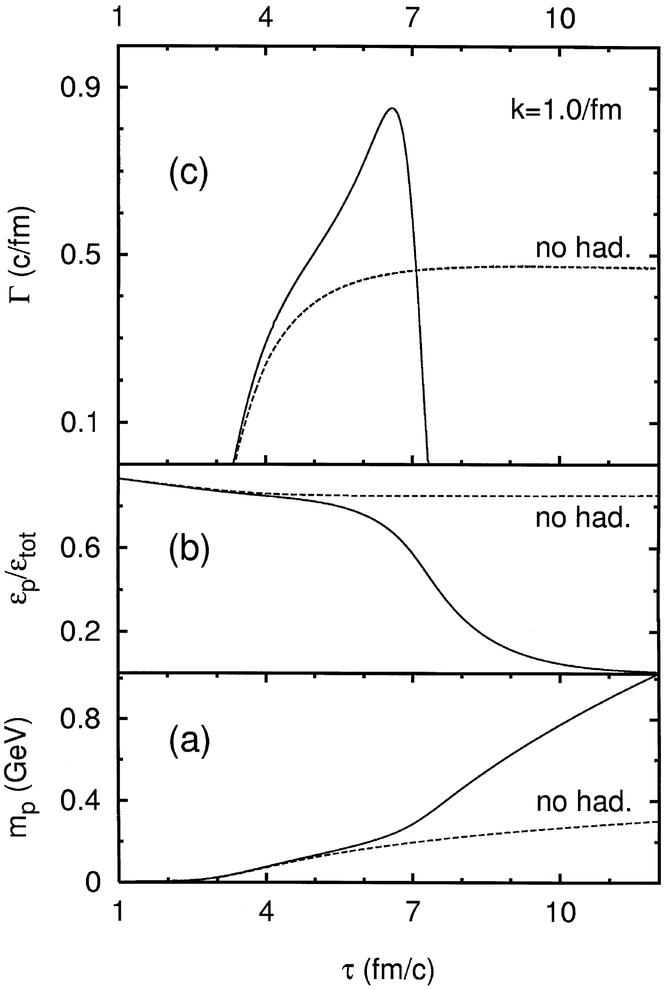

Fig. 1. The time evolution of the parton mass (a), energy ratio (b) and instability growth rate (c) at wave number $k_z = 1/fm$ for the cases with mean field only (dashed lines) and with hadronization (solid lines)

The evolution of the hadronization in the parton plasma can be read from the ratio of the parton energy density to the total energy density,

$$r(\tau) = \frac{\epsilon_p(\tau)}{\epsilon_{tot}(\tau)} = \frac{\epsilon_p(\tau)}{\epsilon_p(\tau)+\epsilon_\pi(\tau)+V(\tau)} , \qquad (14)$$

with the numerical result given in Fig.1b. At the beginning of the evolution, the system is controlled by the expansion and the pion contribution is very small. The slight decrease of the ratio is due to the almost constant potential. When the mass $m_p(\tau)$ approaches the value m_p^c in (11), hadronization becomes the dominating effect. This happens around $\tau = 7fm/c$.

Let us now study the onset of instability associated with the deconfinement phase transition in the expanding parton plasma. In the spirit of linear response theory, we consider a perturbation δf_p around the scaling solution f and study its time dependence. From the Boltzmann equation and the gap equation for the partons, the deviations $\delta f_p(x,\mathbf{p})$ and $\delta m_p(x)$ around the zeroth-order distribution $f_p(x,\mathbf{p})$ and effective parton mass $m_p(x)$ are characterised by the linear equations

$$\partial_t \delta f_p + \mathbf{v}_p \cdot \nabla \delta f_p - \frac{m_p}{E_p} \nabla \delta m_p \cdot \nabla_p f_p = -\delta I_p \ ,$$
$$\left(\frac{\partial^2 V}{\partial m_p^2} + g_p \int \frac{d^3\mathbf{p}}{(2\pi)^3} \frac{\mathbf{p}^2}{E_p^3} f_p \right) \delta m_p$$
$$+ g_p \int \frac{d^3\mathbf{p}}{(2\pi)^3} \frac{m_p}{E_p} \delta f_p = 0 \qquad (15)$$

Here $\mathbf{v}_p = \mathbf{p}/E_p$ is the parton velocity in the plasma, and δI_p is the fluctuation of the loss term,

$$\delta I_p = \frac{\partial}{m_p} (\frac{1}{\theta_p}) f_p \delta m_p + \frac{1}{\theta_p} \delta f_p \ . \qquad (16)$$

In the derivation of the above linear equations we have considered the often used adiabatic approximation in the linear response analysis, namely that compared with the expected strong space-time variation of the fluctuations δf_p and δm_p, the space-time dependence of the zeroth-order quantities f_p and m_p can be neglected. This approximation may be questionable for too fast hadronization.

After a Fourier transform of the linear equations with respect to the perturbations we derive the linear response equation for the wave number $\mathbf{k}$ and frequency ω of the plane waves in our infinite parton plasma,

$$\left[\frac{\partial^2 V}{\partial m_p^2} + g_p \int \frac{d^3\mathbf{p}}{(2\pi)^3} \left(\frac{\mathbf{p}^2}{E_p^3} f_p + \frac{m_p}{E_p} \frac{m_p}{E_p} \mathbf{k} \cdot \nabla f_p \right.\right.$$
$$\left.\left. + \frac{i f_p \frac{\partial}{\partial m_p} (\frac{1}{\theta_p})}{\mathbf{k} \cdot \mathbf{v_p} - \omega - i/\theta_p} \right) \right] \delta m_p(\omega, \mathbf{k}) = 0 \ . \qquad (17)$$

Since we have assumed m_π to be constant, $\delta m_\pi = 0$ and the fluctuation of the pion distribution in the plasma is a consequence of the parton disturbance. This can be seen clearly from the equation

$$\delta f_\pi(\omega, \mathbf{k}, \mathbf{p}) = \frac{i \frac{\partial}{\partial m_p} (\frac{1}{\theta_\pi}) f_\pi}{\omega - \mathbf{k} \cdot \mathbf{v}_\pi - i/\theta_\pi} \delta m_p(\omega, \mathbf{k}, \mathbf{p}) \ . \qquad (18)$$

The fluctuation can propagate in the plasma only if $\delta m_p \neq 0$, hence the response equation which gives the relation between the frequency and the wave number is the zero of the square bracket in (17). Inserting f_p from (13) and let $\mathbf{k} = (0, 0, k_z)$ (perturbation in longitudinal direction) one has

$$\int_0^\infty dp_\perp p_\perp \int_{-\infty}^\infty dp_z \frac{\dot{f}_p^{eq}}{E_p \tilde{E}_p}$$
$$\times \left[\frac{k_z v_z + i \frac{T \tilde{E}_p}{m_p} \left(\frac{\tau_0}{\tau}\right)^2 \frac{\partial}{\partial m_p} \left(\frac{1}{\theta_p}\right) \frac{f_p^{eq}}{\dot{f}_p^{eq}}}{k_z v_z - \omega - i/\theta_p} \right.$$
$$\left. - \frac{d\tau}{dm_p} \left(\frac{p_z^2}{m_p \tau} - \frac{1}{\theta_p} \left(\frac{\tau_0}{\tau}\right)^2 \frac{T \tilde{E}_p}{m_p} \frac{f_p^{eq}}{\dot{f}_p^{eq}} \right) \right]$$
$$= 0, \qquad (19)$$

where $\dot{f}_p^{eq}$ indicates the derivative of the equilibrium distribution with respect to the variable $\frac{\tilde{E}_p}{T}$. The linear dispersion relation in the transverse direction is similar to (19), and the corresponding growth rate is smaller.

Unstable modes are characterised by the imaginary solution of the linear equation (19), $\omega = i\Gamma$. They grow like $e^{\Gamma t}$. In the absence of hadronization, $1/\theta_p = 0$, the fraction containing ω can be divided through by k_z and ω/k_z is the unknown quantity, which leads to $\Gamma \propto k_z$ for the instability rate. This case is discussed in [12], and the instability growth rate, called Γ_{mf}, is shown by the dashed curve in Fig.1c. The system is stable in the early stage of the evolution, but then the instability increases rapidly. With hadronization, θ_p appears at three places in (19). In the denominator, $\omega = i\Gamma$ is changed to $\omega = i(\Gamma + 1/\theta_p)$. From the two other terms, the one in the numerator is more important than the one in the round brackets. We decompose the total growth rate Γ into three parts, the mean-field part, Γ_{mf}, a correction $\delta\Gamma_{hd}$ and the inverse of the relaxation time,

$$\Gamma = \Gamma_{mf} + \delta\Gamma_{hd} - \frac{1}{\theta_p} \ . \qquad (20)$$

The time evolution of the total instability growth rate is indicated by the solid curve in Fig.1c. The small hadronization rate at the early stage does not change the stability of the parton plasma. When hadronization becomes important, first $\delta\Gamma_{hd} > 0$ and it increases the instability rate by nearly a factor two, until finally, the last term in (20) dominates: the growth rate of the instability is over-compensated by the depletion of the density by the overall hadronization.

We want to explain the behaviour of Γ under the influence of the mean field and of hadronization in simple physical terms (Fig.2). We assume a homogeneous density $\rho = constant$ as a zeroth order solution and a fluctuation $\delta\rho$, which is negative in the picture. Via the gap equation, a reduced density leads to an increased mass $\delta m > 0$. Since the Newtonian equation for the quasiparticle model is $\dot{p} = -\nabla E, E^2 = \mathbf{p}^2 + m^2(\mathbf{x}, t)$, a modification in the mass $\delta m > 0$ leads to $\delta E > 0$ and particles will flow away from the edges of the perturbation where the gradient is largest. Therefore $\delta\rho$ will become more negative and we have an instable situation. Since the flow is proportional to $\nabla\delta E \propto \nabla\delta m$, the width for the growth rate is inversely proportional to the size of the perturbation,i. e. $\Gamma_{mf} \propto k_z$. The hadronization rate $1/\theta_p(m_p)$ is a monotonically increasing function of the mass. Therefore $\delta m > 0$ leads to an increased hadronization rate, as shown in the lowest picture of Fig.2. In our model the hadronization is a volume effect (proportional to m_p and not to ∇m_p) and therefore $\delta\Gamma_{hd}$ is not proportional to k_z. If one starts with a positive perturbation $\delta\rho > 0$, one goes though the arguments and finds that $\delta\rho$ increases. The result of the competition between the two self-generating mechanisms and the absorption term $1/\theta_p$ is reflected in the steep growth and then the fast dropping down of the total growth rate Γ in Fig.1c. For the pions, no self-consistency is obtained

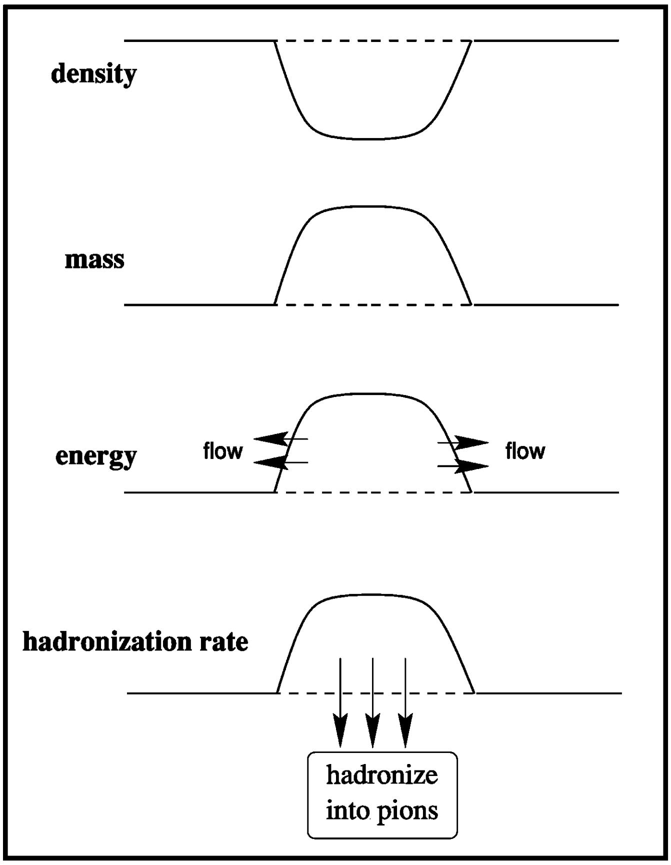

Fig. 2. A schematic picture of the two mechanisms which lead to the self-accelerating growth rate of the instability. From top to bottom: A negative fluctuation of the density, leads to a positive fluctuation of the mass, which generates a "hill" in the total local energy from which the quasiparticles flow away. Finally, an increased mass leads to faster hadronization, again reducing the density

since $\delta m_\pi = 0$. Any fluctuation δf_π is proportional to δm_p and follows the fluctuations of the partons.

In summary, we have investigated the importance of the hadronization on the evolution of the instabilities in an expanding parton plasma by solving the Boltzmann equations with source terms. We have shown that the inclusion of the hadronization effect in the mean-field propagation changes significantly the parton dispersion relation. The unstable modes are corrected by hadronization in two aspects: The transition from partons to pions enhances the instability first, but finally the instability is fully eaten up by the hadronization. While the model contains all essential aspects of QCD phase transition, confinement, chirality and hadronization, and further more the Bjorken scenario for the space time development, the details of the model, quasiparticle approach, no thermalization and a rather schematic expression for the hadronization rate are open for improvement. Yet, qualitatively, the conclusions of the calculation are sound: Hadronization plays a significant role in any understanding of instabilities.

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