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DECONFINEMENT SIGNALS**

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QUARK-HADRON INTERACTIONS IN NUCLEAR MATTER: DECONFINEMENT SIGNALS *

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1 Introduction

As has been noted many years ago, the presence of a phase transition in the equation of state (EoS) may influence in an important way the evolution of a hot and dense system formed in a heavy ion collision. In particular, the EoS near the phase transition becomes 'softer' what may cause some qualitative changes in the hydrodynamic flow pattern. In searching for experimental signals of the deconfinement phase transition, EoS is traditionally treated within a phenomenological two-phase model, i.e. one assumes that a system exists either in the hadronic phase or in the quark-gluon plasma phase. This approach results in the first order phase transition for any system while the QCD lattice calculations clearly demonstrate a large variety of the deconfinement phase transitions depending on system symmetry, flavor and quark content [1].

To avoid discrepancy with the lattice results, our study of EoS is based on the statistical mixed phase model [2]. The underlying physics assumption is that in nuclear

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matter unbound quarks and gluons may coexist with hadron phase forming an *interacting space-homogeneous* quark/gluon-hadron phase. Since the average distance between hadrons and quarks/gluons in this mixed phase may be of the same order as that between hadrons, their interaction with unbound quarks/gluons plays an important role defining in some way the order of the phase transition under question. In the two-phase model, the requirement of thermodynamic consistency leads to the necessity of introducing the so-called Gibbs mixed phase where, in contrast with our treatment, coexisting quark-gluon plasma and hadrons *are spatially separated* and therefore do not interact with each other.

2 Mixed Phase Model

2.1 Model Hamiltonian

In considering the mixed phase, we assume that unbound quarks/gluons and hadrons are good quasiparticles. The separation of quasiparticles implies that instead of working with an exact Hamiltonian, one employs an effective quasiparticle Hamiltonian H which may depend on thermodynamic variables. As a rule, the separation of quasiparticles can not be done exactly, and H should be constructed basing on some physical arguments. In this case, one should take care of thermodynamic consistency because a Hamiltonian depending on thermodynamic variables has to satisfy certain relations which may be called *the conditions of thermodynamic consistency* [2, 3]:

$$\langle \frac{\partial H}{\partial T} \rangle = 0, \quad \langle \frac{\partial H}{\partial \rho_\alpha} \rangle = 0, \quad (1)$$

where $\langle A \rangle$ denotes the average value of the operator A over statistical ensemble and ρ_α is the quasiparticle density for the species α .

To study thermodynamics of the mixed phase we shall use the following effective Hamiltonian [2]:

$$\begin{aligned} H = & \sum_s \int \psi_g^+(\vec{x}, s) \left(\sqrt{-\nabla^2} + U_g(\{\rho\}) \right) \psi_g(\vec{x}, s) d\vec{x} + \\ & + \sum_q \sum_s \int \psi_q^+(\vec{x}, s) \left(\sqrt{-\nabla^2} + m_q^2 + U_q(\{\rho\}) \right) \psi_q(\vec{x}, s) d\vec{x} + \\ & + \sum_{nj} \sum_s \int \psi_{nj}^+(\vec{x}, s) \left(\sqrt{-\nabla^2} + m_{nj}^2 + U_{nj}(\{\rho\}) \right) \psi_{nj}(\vec{x}, s) d\vec{x} - C(\{\rho\}) V, \end{aligned} \quad (2)$$

where $\psi_\alpha(\vec{x}, s)$ denotes a field operator for quasiparticles: $\alpha = g$ means unbound gluons; $\alpha = q$ is unbound quarks ($q = u, \bar{u}, d, \bar{d}, \dots$); $\alpha = nj$ corresponds to n -quark (gluon) hadrons of the j -type. Index s accounts for quantum degrees of

freedom (spin, isospin, color). In (2) $U_\alpha(\{\rho\})$ is the mean field by which the system acts on a quasiparticle of the mass m_α and $\{\rho\}$ is a joint set of individual densities ρ_α of all quasiparticles. Here $C(\{\rho\})V$ is the c -number term (V is the system volume) taking into account the difference between quasiparticle and true vacuum energy. This function can be found by solving eqs.(1) if the mean-field potentials $U_\alpha(\{\rho\})$ are known.

It is worthy to note that both the mean fields and function C in (2) do not depend explicitly on temperature. It is not an extra assumption but the direct consequence of thermodynamic consistency conditions (1) applied to the specific semi-relativistic form of the model Hamiltonian (2) [2, 3, 4].

2.2 Interaction of unbound quarks/gluons with the mixed phase

The mean field affecting color charges in the quark-gluon plasma will depend on the density of surrounding color particles. In the MIT-model of hadrons, the presence of the vacuum pressure corresponds to the action of a mean field inversely proportional to the density of unbound quarks and gluons $\rho_{pl} = \rho_g + \sum_q \rho_q$. On the other hand, if one consider the color charge interaction via the string potential in the nearest-neighbor approximation in analogy with the Debye screening, we arrive at the mean field potential $\sim \rho_{pl}^{-1/3}$. The exact value of the power index is defined by spartial correlations among color particles which are unknown at present beyond the perturbative regime. So, it is reasonable to take the mean field potential of quarks and gluons in plasma in the following form:

$$U(\rho)_q = U(\rho)_g = \frac{A}{\rho^\gamma} \quad (3)$$

with $\rho \equiv \rho_{pl}$ where the parameters A and γ may be found, for example, by fitting the available lattice results on QCD thermodynamics.

One should notice that the assumption (3) mimics two important limiting cases of the QCD interaction. In the limit of $\rho_{pl} \rightarrow 0$ the interaction potential goes to infinity, i.e. an infinite energy should be spent to create an isolated quark or gluon what ensures the confinement of color objects. In other extreme of high energy density (or ρ_{pl}) we arrive at the asymptotic freedom regime.

Up to here, a pure quark-gluon plasma was discussed. How will the interaction of unbound color charges be modified in the mixed phase (MP) due to an admixture of hadrons? Since it is not much known about quark-hadron interaction, it is reasonable to consider two limiting cases. The first one is to use eqs. (3) with $\rho = \rho_{pl}$ for MP as well, which implies neglecting any interaction between hadrons and unbound quarks/gluons. In other words, the presence of free color charges in hadron surrounding does not result in hadron color polarization and these charges are unable 'to see' quarks and gluons confined

in hadrons. The second limit corresponds to (3) with the substitution of *the total density of quarks and gluons* of MP:

$$\rho = \rho_{pl} + \sum_{nj} n \rho_{nj} . \quad (4)$$

This approximation corresponds to a very strong color polarization of hadrons in the mixed phase when free color objects "feel" no difference between bound and unbound quarks/gluons. At first sight, one can hardly make any choice between these two limiting cases because there is no estimate for hadron polarizability in color environment. However, it is easy to see [4] that in the first limiting case neglecting quark-hadron interactions gives rise only to the first order phase transition in contrast with the available lattice data. Thus, as follows from quite general requirements of the confinement for color charges and variety of deconfinement scenario, to describe the unbound quark/gluon interaction with MP, eq. (3) with the total quark/gluon density given by eq. (4) should be used.

Till now, we limited ourselves by the so-called non-perturbative part of the quark-gluon interaction. There is, however, a perturbative part corresponding to the one-gluon exchange interaction. It does not change the order of the phase transition but may manifest itself at the high energy density (temperature) regime. We take into account the perturbative interaction of quarks and gluons by a simple reduction of a number of effective degrees of freedom of unbound quarks and gluons via the following substitution in the free energy density [5]:

$$\xi_q \rightarrow \xi_q \left(1 - \frac{25 g^2}{42 \pi^2} \right) , \quad \xi_g \rightarrow \xi_g \left(1 - \frac{15 g^2}{16 \pi^2} \right) . \quad (5)$$

The running coupling constant g may be taken as a constant value because of its weak temperature dependence. Fitting the lattice data for the purely gluonic $SU(3)$ theory [6] in the region of the QCD phase transition gives $g^2 \approx 10$. This estimate is expected to be valid also in the $SU(3)$ gluon matter with two flavors.

2.3 Interaction of hadrons with the mixed phase

The mean field acting on the nj -hadron in MP may be presented as follows:

$$U_{nj} = \frac{n A}{\rho^\gamma} \left(1 - (1 - w_{pl})^{-\gamma} \right) + U_{nj}^{(h)} . \quad (6)$$

Here the first term is induced by the quark-gluon component of MP and vanishes when there is no unbound quarks/gluons, $w_{pl} \equiv \rho_{pl}/\rho = 0$, w_{pl} being the concentration of quark-gluon plasma in MP considered as an order parameter of the QCD phase transition in this model. The form of this term is coming from the crossing relation $\partial U_{nj}^{(pl)}/\partial \rho_q = \partial U_q/\partial \rho_{nj}$ which, in its turn, is a consequence of conditions of thermodynamic consistency (1) [2].

The remaining term in (6) is the mean field acting on the given hadron in the case of purely hadronic environment, $w_{pl} = 0$, which can be found from ordinary hadronic models. We have extended to our case the non-linear mean-field model proposed by Zimanyi et al [7]

$$U_{nj}^{(h)} = g_{nj}^r \varphi_1(x) + g_{nj}^a \varphi_2(y) , \quad (7)$$

with

$$x = \sum_{mi} g_{mi}^r \rho_{mi}, \quad y = \sum_{mi} g_{mi}^a \rho_{mi}$$

where the repulsive $g_{mi}^r > 0$ and attractive $g_{mi}^a < 0$ coupling constants are assumed to scale with respect to a number n of constituent quarks:

$$g_{nj}^r = n g_r, \quad g_{nj}^a = n g_a . \quad (8)$$

Functions $\varphi_1(x)$ and $\varphi_2(y)$ are defined by the equations [7]:

$$b_1 \varphi_1 = x, \quad -b_1(\varphi_2 + b_2 \varphi_2^3) = y \quad (9)$$

with free constants b_1 and b_2 . After some algebra one can find that this scheme is really ruled by only three free parameters:

$$c_1 = \frac{b_1}{g_r^2}, \quad c_2 = \frac{b_1}{g_a^2}, \quad c_3 = \frac{b_1 b_2}{g_a^4}$$

which values are fixed by properties of the ground state ($T = 0$ and $n_B = n_0 \approx 0.17 \text{ fm}^{-3}$) of nuclear matter: Pressure equals to zero, binding energy is -16 MeV and compressibility is 210 MeV (see details in [4]). As is shown in Fig. 1, thermodynamic properties of a non-ideal pion gas are successfully described within the scaling hypothesis (8) in the close agreement with results of modern hadronic model [8] (for more discussion see [4]).

One should note that for gluonic clusters, glueballs, the mean field $U_{nj}^{(h)}$ is not so well known as for hadrons and it was taken in the Hartree approximation as follows:

$$U_{nj}^{(h)} = \sum_{mi} \Phi_{nj,mi} \rho_{mi} \quad (10)$$

with $\Phi_{nj,mi} = \int \Phi_{nj,mi}(r) d\vec{r}$. In this situation the scaling relation (8) for the interaction potential between the nj -th and mi -th clusters can be represented as $\Phi_{nj,mi}(r) = n m \Phi(r)$. So, to use (10) one needs an additional free constant to characterize the glueball-glueball interaction strength [2, 3].

As to A and γ parameters, the structure of the MP model does not leave much room for variation of their values. Fitting the lattice data on the QCD thermodynamics for $SU(2)$ and $SU(3)$ quarkless systems [2, 3] gives a good agreement between the MP model

Figure 1: The temperature dependence of energy density (left panel) and pressure (right panel) for a pion gas within the statistical MP model. Dashed and dotted curves are the Rapp-Wambach results [8] for interacting and ideal pions, respectively.

and the lattice results at $\gamma \approx 0.6 \div 0.65$ for the $SU(2)$ symmetry and at $\gamma \approx 0.5 \div 0.65$ for the $SU(3)$ color group, being the best at $\gamma \approx 0.62$ in both the cases. Below, the parameter γ is fixed just by this value. Describing the intensity of color interaction, the parameter A is defined by the QCD running constant α_S . The passage from the quarkless $SU(3)$ case to that of the $SU(3)$ theory with two light flavors is accompanied by an increase in α_S by about 10%. So, one may expect to describe both these cases by the same value of A . The MP model analysis of thermodynamics of the $SU(3)$ gluonic system showed that $A^{1/(3\gamma+1)} \simeq T_{dec}$ [2]. According to the recent lattice results [9], the deconfinement temperature of the quarkless $SU(3)$ system is estimated as $T_{dec} = 260 \text{ MeV}$. Therefore, for the MP model parameters γ and A we have:

$$\gamma = 0.62, \quad A^{1/(3\gamma+1)} \simeq 260 \text{ MeV} . \quad (11)$$

Being limited to the case of two light quark flavors, the unbound u and d quarks together with their antiquarks and unbound gluons are taken into consideration. The current quark mass m_q is 7 MeV . In the hadronic sector, all mesons up to $\phi(1020)$ made of u , d quarks/antiquarks are taken into account. Besides baryons, nucleons and Δ isobars, their antiparticles are included, as well. This set of hadrons seems to be sufficient to get necessary accuracy for general thermodynamic quantities.

3 Mixed Phase Thermodynamics

3.1 Comparison with the lattice QCD results

Since the Hamiltonian (2) is given in the mean field approximation, the free energy is calculable exactly

$$f = \sum_{\alpha} \frac{T \xi_{\alpha} \nu_{\alpha}}{2\pi^2} \int_0^{+\infty} \ln \left[1 - \nu_{\alpha} \exp \left(- \frac{\sqrt{k^2 + m_{\alpha}^2} + U_{\alpha} - \mu_{\alpha}}{T} \right) \right] k^2 dk + \sum_{\alpha} \mu_{\alpha} \rho_{\alpha} - C(\rho), \quad (12)$$

where μ_{α} is the chemical potential and ν_{α} takes the value 1 for bosons and -1 for fermions. Then, all other thermodynamic quantities may be easily obtained by the proper differentiation of the free energy (12).

Figure 2: Temperature dependence of relative gluon abundance (left panel) and non-ideality (right panel) of pure gluonic $SU(3)$ system. The lattice data are taken from [9].

The comparison with the lattice results for pure gluonic $SU(3)$ case is presented in Fig. 2. The two lowest glueball states $M_G(0^{++}) = 1470 \text{ MeV}$ and $M_G(2^{++}) = 2205 \text{ MeV}$ estimated in the lattice calculations [10] are taken into account. As is seen, the order parameter $W_g = \rho_g/(\rho_g + 2\rho_G)$ exhibits the clear first order phase transition with a characteristic 'hysteresis' loop. The deconfinement temperature $T_{dec} = 260 \text{ MeV}$ predicted by the lattice results [9] is followed from the choice $A^{1/(3\gamma+1)} = 250 \text{ MeV}$. This value of the parameter A will be used below in the case of $SU(3)$ system with two light flavors. One should note that admixture of unbound gluons is negligible at temperatures below T_{dec} (see the insert in Fig. 2) but the glueball concentration amounts 20 – 10% at $T > T_{dec}$

resulting together with strong interaction of unbound gluons in non-ideal thermodynamic behavior at high temperatures. The last point is demonstrated in the right-hand panel of Fig. 2 by the direct comparison with the lattice data for the temperature dependence of $(\epsilon - 3p)$ which may be treated as a measure of non-ideality.

Figure 3: The reduced energy density and pressure of the $SU(3)$ system with two light flavors at $n_B = 0$ calculated within the MP model for $g^2 = 8$ (curve 1) and $g^2 = 9$ (curve 2). Triangles and squares are the lattice QCD results obtained in the Wilson [11] and Kogut-Susskind [12] schemes, respectively.

The case of $SU(3)$ theory with two flavor quarks at $n_B = 0$ is given in Fig. 3. Presented for two schemes of discretization of the QCD action, the Wilson scheme [11] and Kogut-Susskind [12] one, the lattice data agree with each other in both the crossover type transition and the value of the deconfinement temperature: The first scheme gives $T_{dec} \approx 150 \text{ MeV}$ [1]; the second one is $T_{dec} = 143 \div 154 \text{ MeV}$ [12]). For the above cited parameters the MP model predicts $T_{dec} = 153 \text{ MeV}$ to be determined by the position of a maximum of heat capacity. One should note that the Kogut-Susskind scheme has a problem with the description of hadrons: The energy density and pressure go to zero too fast as compared with that expected for a pion gas at $T < T_{dec}$ and $n_B = 0$. This is possibly coming from the point that a number of temporal steps in these calculations is not large enough and due to that the ratio of the calculated pion mass to the temperature of the QCD phase transition is still too high, $m_\pi/T_{dec} = 1.94$, though the used values of the light quark masses are close to the physical one: $m_q/T_{dec} = 0.075$.

3.2 Softest point effect in the MP model

Shown above in Fig. 3, EoS for the $SU(3)$ gluonic matter with quarks of two light flavors is replotted in left panel of Fig. 4 but in the other representation: p/ϵ ratio vs ϵ . The lattice data in the Wilson scheme and the MP model results exhibit a minimum in the p/ϵ function in the region of the deconfinement transition predicting the softest point (SP) position as large as $\epsilon_{SP} \approx 0.5 \text{ GeV}/\text{fm}^3$. The recent results obtained in the Kogut-Susskind discretization scheme [12] plotted in Fig. 4 do not unfortunately allow one to make any definite conclusion as to the SP existence and its position due to the above-mentioned problem resulting in a strong underestimation of the pressure at $T < T_{dec}$. In the same figure one can see also the two-phase model data with the bag constant $B = 235^4 \text{ MeV}^4$ used. This value of the vacuum pressure B results in a reasonable value of $T_{dec} \approx 160 \text{ MeV}$ but gives a too large value for SP $\epsilon_{SP} \approx 1.5 \text{ GeV}/\text{fm}^3$ to be hardly conciliated with the lattice data predictions.

Figure 4: The p/ϵ -representation of EoS for the two flavor $SU(3)$ system at different baryon densities n_B . Triangles and squares are the lattice QCD results obtained in the Wilson [11] and Kogut-Susskind [12] schemes, respectively. The lines are calculated within the MP (solid), two-phase (dashed) and hadronic gas (dotted) models.

As is seen from the right panel Fig. 4, the baryon density dependence of the p/ϵ function in the MP model differs essentially from the two-phase model predictions: with increasing n_B the function minimum is washed out and disappears completely at $n_B \gtrsim 0.4 n_0$. This leads to quite controversial conditions to observe the SP effect in heavy-ion collisions. Indeed, being specific for the crossover transition, the low value of ϵ_{SP} testifies that the formation of a long-lived fireball as a signal of the deconfinement phase

transition might be observed at rather moderate energies, below 10 $A\text{GeV}$. However, due to high stopping at these energies, the fireball formed will have a high baryon density which, according to the MP model but in contrast with the two-phase model, should result in washing out the minimum in p/ϵ and degenerating the SP effect. It is noteworthy that within the MP model EoS for baryon-rich matter turns out to be closer to that for interacting hadronic gas. Some alternative case is given by heavy-ion collisions at ultra-relativistic energies of LHC and RHIC colliders. If the Bjorken mechanism [15] is realized at these energies, the manifestation of the SP effect is possible in the course of evolution of a baryon-free fireball.

To see experimental consequences of the mixed phase EoS, the latter should be incorporated in a hydrodynamic-type model. In the first approximation, we consider an expanding blob of the compressed and heated QCD matter (a fireball) formed in heavy-ion collisions. The initial state for this fireball is defined within the quark-gluon string model [16]. Further fireball evolution is described as an isentropic scaled expansion. According to this prescription, the volume densities of the conserved quantities, entropy and baryon number, are evolved inversely proportionally to the expansion time; so all the other thermodynamic quantities can now be easily found for the given EoS at every temporal step of evolution.

Figure 5: The correlation between p/ϵ and the energy density reached during the expansion of a fireball formed in central $Au + Au$ collisions at E_{lab} . The results are obtained within the MP (the left-hand side) and two-phase (the right-hand side) models.

In Fig. 5 the key evolution quantity p/ϵ is shown as a function of the energy density attained during expansion. The absence of a long-lived fireball in the MP model is a direct consequence of the mixed phase EoS where the softest point is washed out with an increase

in n_B , as noted above. Making use of the two-phase model essentially overestimates the SP effect resulting in noticeably longer life-time of a fireball.

The same evolution for both models is shown in Fig. 6 in the $T - \mu_B$ phase plane. Though the initial values of ε and n_B are the same, the modelling paths are drastically different but their final points are rather close each to other. This means that these models are nor easily distinguishable in hadronic observables but the difference in trajectories may be probed by leptons.

Figure 6: Dynamical trajectories in $T - \mu_B$ plane for evolution of a fireball formed in central $Au + Au$ collisions at E_{lab} . Full and open symbols correspond to the mixed phase and two-phase model results, respectively. Circles with error bars are experimental points from [18,19].

Recently [17, 18, 19], temperature and chemical potential were estimated experimentally at 200, 10 and 2 AGeV for the chemical freeze-out point with using the two-phase model. These results are plotted in Fig. 6 too. It is noteworthy that Particularities of the statistical mixed phase model for describing the deconfinement phase transition and physics constraints on the model parameters, are discussed. The modifications proposed concern an improved treatment of hadron-hadron interactions within the nonlinear mean-field model of nuclear matter and an inclusion of the one-gluon exchange correction into the quark-gluon sector of the mixed phase thermodynamic potential. It is shown that the mixed phase model successfully reproduces both the available lattice data on the QCD deconfinement transition and thermodynamic characteristics of hadron systems evaluated within the modern models. The 'softest point' of the equation of state resulting in a possible formation of a long-lived fireball in nuclear collisions (the 'softest point' effect) is analyzed. The obtained results are confronted with the data on the two-phase model.

Some manifestation of the mixed phase formation and deconfinement transition is estimated in fireball expansion dynamics. We discuss experimental possibilities of observing the softest point effect as a signal of the deconfinement phase transition. These experimental values turn out to correlate with the turning point of the dynamical path for the MP model, being below the phase boundary for the two-phase model at moderate colliding energies.

4 Concluding remarks

So, the realized hypothesis on coexistence of interacting hadrons and unbound quarks/gluons in nuclear matter allows one to reproduce successfully the lattice QCD results as to both the critical temperature and the order of a phase transition for different systems as well as to extend the developed mixed phase model to baryon-rich systems in a parameter-free way. The MP model predicts that the softest point is located at comparatively low values of the energy density: $\varepsilon_{SP} \approx 0.5 \text{ GeV}/fm^3$. This value of ε is close to the energy density inside a nucleon and thus reaching this value signals us that we are dealing with a single 'large' hadron consisting of deconfined matter. For baryonic matter the softest point is washed out at $n_B \gtrsim 0.4 n_0$. As an experimental consequence, an excitation function for the directed nucleon flow within the MP model is expected to behave quite similarly to the case of interacting hadron gas. However, this expectation is essentially different from the two-phase model of the QCD deconfinement predicting a deep minimum in the flow excitation function at $E_{lab} \approx 5 \text{ AGeV}$ [13].

The directed flow is certainly not the only signal of the mixed phase and deconfinement phase transition. A generalization of the MP model to the case of three flavors [20] shows a sharp increase in the relative fraction of strange quarks to be observed at $\varepsilon \gtrsim \varepsilon_{SP}$, i.e. the strangeness enhancement is correlated with the SP position. The available systematics of the strange particle production [21] gives some evidence in favor of such a behavior: the reduced strangeness suffers a sharp increase somewhere at $E_{lab} = 5 - 10 \text{ AGeV}$ and then is practically independent of the colliding energy. The identical particle interference deserves a particular attention since it is sensitive to 'granularity' of an emission source [22]. Alongside with the source radius, the correlation length for unbound quarks/gluons will play a role of a new additional scale in the MP case. The production of dilepton pairs and hard photons is of great interest. The specific path of a system in the $T - \mu_B$ space corresponding to the given EoS will result in emission probabilities different for different EoS. In addition, taking into account interactions between hadrons and unbound quarks/gluons opens a new channel for dilepton and soft photon production. In particular, this channel may contribute to the enhanced yield of dileptons with the effective masses $0.2 - 0.6 \text{ GeV}$

recently observed by the CERES Collaboration in relativistic heavy-ion collisions [23]. One should note also that the presence of unbound gluons in MP allows one to speculate as to an origin of anomalous J/Ψ suppression in the latest $Pb + Pb$ data at the CERN energies [24].

The discussed consequences of a possible manifestation of the MP formation in heavy-ion collisions are of qualitative nature. For quantitative predictions the developed mixed phase EoS should be included into full hydrodynamic calculations. This work is in progress now.

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