Cavitation and bubble collapse in hot asymmetric nuclear matter.

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Abstract

The dynamics of embryonic bubbles in overheated, viscous and non-Markovian nuclear matter is studied. We show that the memory and the Fermi surface distortions significantly affect the hinderance of bubble collapse and determine a characteristic oscillations of the bubble radius. These oscillations occur due to the additional elastic force induced by the memory integral.

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1. Growing or collapsing bubbles in an overheated (undercompressed) fluid has been an intriguing problem of physics over a long period of time \[1\]. In nuclear physics, the bubbles may be formed in the expanded nuclear matter in a heavy-ion reaction then providing the break-up of the matter into fragments \[2, 3, 4\]. The formation of embryonic bubbles in hot nuclear matter due to the quantum and statistical fluctuations has been studied earlier in Refs. \[5, 6, 7, 8, 9, 10\]. The goal of the present paper is the analysis of the dynamic evolution and the collapse of the bubble in an overheated nuclear matter.

In a classical liquid, the dynamics of the bubble of radius $R$ is derived mainly from the thermodynamic potential $\Phi(R)$, which reaches a maximum value for the critical radius $R = R^*$ \[12\]. More complicated bubble dynamics occurs in the case of a Fermi liquid where the dynamic distortion of the Fermi surface produces an additional pressure tensor \[13, 14\]. Moreover, the Fermi surface distortion and the interparticle collisions lead here to the non-Markovian equations of motion for the relevant collective variables \[15\]. Below the non-Markovian dynamics will be applied to the problem of the collapse of bubbles. The non-Markovian equation of motion for the bubble radius $R$ will be derived in a hot nuclear Fermi-liquid. However, the final macroscopic equations of motion including the memory kernel can be applied to a wide number of non-Newtonian liquids.

2. We will consider the dynamics of the bubble (cavitation) with an arbitrary undercritical size $R \leq R^*$, where the critical radius $R^*$ of the bubble in an overheated liquid at the temperature $T$. The critical radius $R^*$ depends on the overheating temperature $\Delta T = T - T_0$, where $T_0$ is the boiling temperature in the case of plane geometry. For fixed values of temperature $T$ and liquid phase pressure $P_{\text{liq}}$ the critical radius $R^*$ is derived from the condition of thermodynamic equilibrium of the vapor bubble with the surrounding liquid and it is given by \[11, 12\]

$$R^* = \frac{2\sigma T_0}{\rho_{\text{vap}} \phi \Delta T}.$$  \hspace{1cm} (1)

Here, $\sigma$ is the surface tension coefficient, $\rho_{\text{vap}}$ is the particle density in the bubble and $\phi$ is the latent heat of evaporation. Considering the dynamics of the surrounding Fermi liquid, one can reduce the collisional kinetic equation for the phase-space distribution function to Euler-like equations for the velocity field $\mathbf{u}(\mathbf{r}, t)$ in the form (for details, see Refs. \[14, 15, 16\])

$$m \frac{\partial}{\partial t} u_\nu + m \rho (u_\mu \nabla_\mu) u_\nu + \frac{1}{\rho} \nabla_\nu P + \nabla_\nu \frac{\delta \epsilon_{\text{pot}}}{\delta \rho} = -\frac{1}{\rho} \nabla_\mu P'_{\nu\mu} + F_{\nu,\text{ext}},$$  \hspace{1cm} (2)
where $\rho \equiv \rho(r, t)$ is the particle density, $\epsilon_{\text{pot}} \equiv \epsilon_{\text{pot}}(r, t)$ is the potential energy density and $\mathcal{P} \equiv \mathcal{P}(r, t)$ is the pressure caused by both the thermal and Fermi motions of nucleons. The pressure tensor $\mathcal{P}_{\nu\mu}' \equiv \mathcal{P}_{\nu\mu}'(r, t)$ in Eq. (2) is caused by the Fermi surface distortion. In the case of the most important quadrupole distortion of the Fermi surface, the pressure tensor $\mathcal{P}_{\nu\mu}'$ satisfies the following equation \cite{15}

$$
\frac{\partial}{\partial t} \mathcal{P}_{\nu\mu}' + \mathcal{P} \frac{\partial}{\partial t} \Lambda_{\nu\mu} = -\frac{1}{\tau} \mathcal{P}_{\nu\mu}' \quad \text{with} \quad \Lambda_{\nu\mu} = \nabla_{\nu}\chi_{\mu} + \nabla_{\mu}\chi_{\nu} - \frac{2}{3} \delta_{\nu\mu} \nabla\chi \chi,
$$

(3)

where $\tau$ is the relaxation time due to interparticle collisions and $\chi_{\nu} \equiv \chi_{\nu}(r, t)$ is the displacement field which is related to the velocity field by $u_{\nu}(r, t) = \partial \chi_{\nu}(r, t)/\partial t$. In the case of a spherical bubble, the displacement field is given by \cite{7, 17}

$$
\chi_{\nu}(r, t) = R^{2}r_{\nu}/3r^{3} \quad \text{for} \quad r \geq R.
$$

(4)

We point out the quadrupole distortion of the Fermi surface does not distort the spherical shape of the bubble if the displacement field $\chi_{\nu}(r, t)$ is given by Eq. (4).

The external force $F_{\text{ext}}$ in Eq. (2) is caused by the vapor pressure on the liquid. The external force on the bubble surface in the radially outward direction per unit area is given by \cite{18, 19}

$$
\bar{F}_{\text{ext}} = \mathcal{P}_{\nu\mu}'(r, t) - \mathcal{P}_{\nu\mu}^0,
$$

(5)

where $\mathcal{P}_{\nu\mu}$ is the vapor pressure within the bubble and $\mathcal{P}_{\nu\mu}^0$ is the pressure of the saturated vapor with respect to a plane surface. Note that the condition of equilibrium between the bubble vapor and the liquid is expressed by $\mathcal{P}_{\nu\mu}^0 = \mathcal{P}_{\nu\mu}'$ and $\mathcal{F}_{\text{ext}} = 0$, see Ref. \cite{12}.

The Euler-like equations (2) can be reduced to the equation of motion for the radius $R(t)$ of the bubble. Multiplying Eq. (2) by $\chi_{\nu}(r, t) = R^{2}r_{\nu}/r^{3}$, summing over $\nu$, integrating over $r$-space and using Eqs. (3), we obtain the non-Markovian equation for the variable $R(t)$

$$
B \ddot{R} + \frac{1}{2} \frac{\partial B}{\partial R} \dot{R}^2 + I(R; t) = -\frac{\partial \Phi}{\partial R} - 4\pi R^2 \mathcal{P}_{\nu\mu}' \left(1 - \frac{\mathcal{P}_{\nu\mu}'}{\mathcal{P}_{\nu\mu}^0}\right),
$$

(6)

where

$$
I(R; t) = B \int_{t_0}^{t} dt' R(t') \exp[(t - t')/\tau] \mathcal{K}(t, t')
$$

(7)

is the memory integral. The inertial parameter $B \equiv B(R)$ and the memory kernel $\mathcal{K}(t, t')$ in Eqs. (6) and (7) are given by \cite{16, 19, 20}

$$
B(R) = 4\pi m \rho_0 R^3, \quad \mathcal{K}(t, t') = \frac{8 \epsilon_F}{5m} \frac{R(t)R(t')}{R(t)R(t')},
$$

(8)
FIG. 1: Dependence of the thermodynamical potential $\Delta \Phi$ (see Eq. (10)) of metastable (overheated) asymmetric nuclear matter on the radius of the bubble.

where $\rho_0$ is the bulk density of the nuclear matter and $\epsilon_F$ is the Fermi energy. Denoting the value of thermodynamic potential $\Phi$ corresponding to the absence of the bubble by $\Phi_0$, we have

$$\Phi \equiv \Phi(R) = \Phi_0 + \Delta \Phi(R), \quad (9)$$

where $\Delta \Phi(R)$ is the change of the thermodynamic potential due to the variation of the bubble radius. The value of $\Delta \Phi(R)$ is given by [12]

$$\Delta \Phi(R) = 4\pi \sigma \left( R^2 - \frac{2R^3}{3R^*} \right). \quad (10)$$

The thermodynamic potential $\Phi(R)$ for $R = R^*$ does not have a minimum value, as in the case of equilibrium, but a maximum value, see Fig. 1. The thermodynamic potential of Eq. (10) corresponds to the metastable liquid phase with $\mu_{\text{liq}} > \mu_{\text{vap}}$, where $\mu_{\text{liq}}$ and $\mu_{\text{vap}}$ are the chemical potentials for the liquid and the bubble vapor, respectively.

Note that in the Markovian limit at $\tau \to 0$ the equation of motion (6) is transformed to the generalized Rayleigh-Plesset equation of classical bubble dynamics [19].
The critical radius $R^*$ in Eq. (10) depends on the temperature $T_0$ of liquid-gas phase transition, the overheating temperature $\Delta T$ and the isotopic asymmetry parameter $X_{\text{liq}}$ ($X = (\rho_n - \rho_p)/(\rho_n + \rho_p)$, where $\rho_n$ and $\rho_p$ are the neutron and proton densities, respectively), see Eq. (11) and Refs. [12, 20]. To evaluate the critical radius $R^*$ we point out that the caloric curve has the plateau region at the temperature $T_0$ [21]. Below we will adopt $T_0 = 4$ MeV and $X_{\text{liq}} = 0.1$. If one assumes the process of isobaric heating for the description of the plateau region in the caloric curve at $T_0 = 4$ MeV, the order of magnitude of the pressure $P_0$ should be $P_0 \approx 10^{-3}$ MeV/fm$^3$ for this process [21]. Under these conditions, one obtains from Eq. (11) $R^* = q/\Delta T$ with $q \approx 4.1$ MeV · fm for a symmetric nuclear matter. The parameter $q$ grows slightly with the asymmetry of the nuclear matter.

In the case of boiling of asymmetric nuclear matter, the vapor asymmetry $X_{\text{vap}}$ significantly exceeds the corresponding liquid asymmetry $X_{\text{liq}}$ (at $X_{\text{liq}} > 0$), see Ref. [21]. This is a feature of the nuclear matter where the structure of the isospin symmetry energy provides the condition $|\mu_n| < |\mu_p|$ and the preferable evaporation of neutrons. Using this fact, we will neglect the proton fraction in the bubble vapor. The presence of the noncondensable neutron vapor within the bubble then gives a contribution to the term $\sim (1 - P_{\text{vap}}/P_0)$ in Eq. (6). (Note that the analogous term in the classical Rayleigh-Plesset equation is due to the assumption that the bubble contains some quantity of contaminant noncondensable gas [19].) The vapor pressure in the bubble includes both quantum and thermal contributions. Since the temperature region of interest is $T \ll \epsilon_F$, the thermal pressure is relatively small and for the degenerate neutron gas one can use the Thomas-Fermi approximation with $P_{\text{vap}}/P_0 = (R^*/R)^5$ in Eq. (6).

The surface tension coefficient $\sigma$ in Eq. (10) is temperature dependent. We will use the following expression for $\sigma \equiv \sigma(T)$ [22]

$$\sigma(T) = \sigma(0) \left[\left(T^2_{\text{crit}} - T^2\right)/\left(T^2_{\text{crit}} + T^2\right)\right]^{5/4},$$

(11)

where $T_{\text{crit}} = 14.6$ MeV is the critical temperature for infinite nuclear Fermi-liquid, associated with the SKM interaction [21].

3. Below we will consider the collapse phase, i.e. the descent from the top of the barrier in Fig. 1 toward $R < R^*$. In general, an accurate evaluation of the bubble dynamics requires to include a thermal boundary condition at the bubble wall to provide the possible changes of the pressure $P_{\text{vap}}$ of the bubble vapor and its temperature [19]. We will neglect the thermal
effects, assuming that the main contribution to the pressure of the bubble vapor is due to the quantum motion of nucleons because of $T \ll \epsilon_F$. If the collapse process is fast enough and the condensation of the neutron’s vapor within the bubble is negligible, the force $\overline{F}_{\text{ext}}$ of Eq. (5) has then to be taken into consideration in Eq. (2). In an opposite case of a slow collapse process, the possible growth of the pressure $P_{\text{vap}}$ at the decrease of the bubble radius $R(t)$ is compensated by the condensation of the vapor providing that $P_{\text{vap}} = P_0 = \text{const}$ and $\overline{F}_{\text{ext}} = 0$.

The non-Markovian equation (6) can be solved if one rewrites it as a set of two connected equations for the bubble radius $R(t)$ and the memory integral $I(R; t)$ with the relevant initial conditions. The bubble starts to collapse from the metastable state (point A on the top of barrier in Fig. 1) at $R(0) = R^*$. Taking the overheating temperature $\Delta T = 0.5$ MeV, we obtain from Eq. (1) $R(0) = R^* = 8.2$ fm. This critical radius $R^* = 8.2$ fm seems too large for the finite nuclei. We point out, however, that the inclusion of the Coulomb forces, which is ignored in our consideration of the infinite nuclear matter, decreases the value of $R^*$ [6, 7]. Note, however, that our consideration cannot be extended to very small values of $R^*$ because for a small enough bubble the finite size effects, in particular the finite diffuse layer for the vapor density inside the bubble, can be important [7]. Using the above mentioned initial condition for $R(0)$, we restrict ourselves to an infinite nuclear matter and expect that the main conclusions concerning the Fermi motion effects on the collapse of the bubble in hot Fermi liquid can be also applied to finite nuclei.

The initial velocity $\dot{R}(0)$ can be derived using the initial kinetic energy $E_{\text{kin,0}}$. Assuming the equipartition of energy over degrees of freedom at $R = R^*$, we use $E_{\text{kin,0}} = (T_0 + \Delta T)/2$ and obtain $\dot{R}(0) = -\sqrt{2E_{\text{kin,0}}/B(R^*)} = -2.9 \cdot 10^{-3}c$, where sign ”−” is used because of the collapse process with $R(t) < R^*$.

In Fig. 2, we show the time dependence of the radius $R(t)$ in the presence of the noncondensable vapor of neutrons with $\overline{F}_{\text{ext}} \neq 0$ (solid lines). Here we have adopted $\epsilon_F = 37$ MeV, $\sigma(0) = 1.1$ MeV $\cdot$ fm$^{-2}$ and $\rho_0 = 0.16$ fm$^{-3}$. For the relaxation time $\tau$ we have used $\tau = \tau_0/T^2$ with $\tau_0 = 850$ fm/$c$ [23]. As seen from Fig. 2, the collapse of the bubble is accompanied by oscillatory behavior of its radius $R(t)$. These oscillations occur due to the memory integral in Eq. (6) and disappear in the limit of Markovian motion (see dashed line in Fig. 2). As seen from Fig. 2, the collapse process stops at a certain values of the bubble radius, $R_{\text{min}} = 0.28 R^*$. The presence of non-zero $R_{\text{min}} \neq 0$ is due to the pressure of the noncond-
FIG. 2: Dependence of the radius $R(t)$ of the collapsing bubble on the time obtained for two regimes: non-Markovian motion, Eq. (6), (solid line) and Markovian (no memory) motion, Eq. (13), with the friction coefficient $\gamma$ from Eq. (14) (dashed line). The inset shows the collapse process near the top of barrier of $\Phi(R)$. The values of $T_0$, $\Delta T$ and $X_{liq}$ as in Fig. 1.

densable neutron vapor in the bubble. The value of $R_{\text{min}}$ is derived by the condition of the compensation of the adiabatic forces in Eq. (5) at $t \gg \tau$.

The memory integral in Eq. (3) leads to a significant delay in the collapse process. To show this effect we will derive the friction coefficient $\gamma$ from the memory integral $I(R; t)$. In a immediate vicinity of the critical radius $R^*$, the solution to Eq. (3) takes the form of a superposition of exponential functions $\exp(\lambda_i t)$ ($i = 1, 2$ and $3$) with eigenvalues $\lambda_i$ obtained as solutions to the secular equation

$$
(\lambda^2 + k/B^*)(\lambda + 1/\tau) + (\tilde{\kappa}/B^*)\lambda = 0.
$$

(12)

Here $k = -4\pi(2\sigma - 5P_0R^*)$, $\tilde{\kappa} = (32/5)\pi \rho_0 \epsilon_F R^*$ and the mass coefficient $B^* \equiv B(R^*)$ is taken from Eq. (8) at $R = R^*$. In the case of the zero-relaxation-time limit, $\tau \to 0$, one obtains from Eq. (12) the motion with $\lambda = \pm \sqrt{|k|/B^*}$, i.e., the time evolution is derived by the static stiffness coefficients $k$. In the opposite case of rare collisions, $\tau \to \infty$, the
solution to Eq. \[12\] leads to a motion with \( \lambda = \pm i \sqrt{(-|k| + \tilde{\kappa})/B^*} \), where the additional contribution to the stiffness coefficient, \( \tilde{\kappa} \), appears because of Fermi surface distortion effect. In both limits \( \tau \to 0 \) and \( \tau \to \infty \), the macroscopic equation of motion \[6\] is reduced to the usual Markovian (no memory) equation with a friction coefficient \( \gamma \)

\[
m\rho_0 R \ddot{R} + \frac{3}{2} m \rho_0 \dot{R}^2 + \gamma \dot{R} = -\frac{2\sigma}{R} \left( 1 - \frac{R}{R^*} \right) - P_0 \left( 1 - \left( \frac{R^*}{R} \right)^5 \right). \tag{13}
\]

The friction coefficient \( \gamma \) is related to the relaxation time \( \tau \) as \[15\]

\[
\gamma = \omega_F B^* \frac{\omega_F \tau}{1 + (\omega_F \tau)^2}, \tag{14}
\]

where \( \omega_F = \sqrt{\tilde{\kappa}/B^*} \) is the characteristic frequency for the eigenvibrations caused by the Fermi surface distortion. The corresponding solution for \( R(t) \) obtained from the Markovian equation of motion \[13\] with friction coefficient \( \gamma \), from Eq. \[14\], is shown in Fig. 2 by the dashed lines. As seen from Fig. 2 the presence of memory effect (solid line) strongly hinders the collapse process. If the collapse process becomes slow enough the growth of the vapor pressure \( P_{\text{vap}} \) is compensated by the condensation of the vapor and the contribution from \( \mathcal{F}_{\text{ext}} \) is negligible in Eq. \[6\]. The collapse process then leads to the disappearance of the bubble.

In Fig. 3 we have plotted the dependence of the time, \( t_{\text{clps}} \), required for the total collapse from \( R = R^* \) to \( R = 0 \), on the relaxation time \( \tau \), for the bubble in the absence of the force \( \mathcal{F}_{\text{ext}} \) (i.e., neglecting the term \( \sim P_0 \) in Eqs. \[6\] and \[13\]).

In Fig. 3 we compare the results for \( t_{\text{clps}} \) for both the non-Markovian motion derived by Eq. \[6\] (solid line) and the usual Markovian (no memory) one given by Eq. \[13\] (dashed line). As seen from Fig. 3, the behavior of \( t_{\text{clps}} \) is changed significantly due to the memory effects. A delay in the collapse of the bubble occurs due the non-Markovian (memory) effect for the large values of the relaxation time.

In the limit of a non-viscous liquid, \( \tau \to 0 \), we obtain for both motions \( t_{\text{clps,0}} = 4 \cdot 10^{-22} \) s. This result can be compared with the classical one of Rayleigh \( t_{\text{clps,R}} \) \[24\]. To do that, we note that under the condition used above for \( T_0 = 4 \) MeV, \( P_0 \approx 10^{-3} \) MeV/fm\(^3\), \( X_{\text{liq}} = 0.1 \) and \( R^* = 8.2 \) fm, the Fermi liquid is undercompressed by a pressure \( \Delta P \approx 0.25 \) MeV/fm\(^3\). The Rayleigh collapsing time \( t_{\text{clps,R}} \) is given by \[19\]

\[
t_{\text{clps,R}} = 0.915 \left( \frac{m \rho_{\text{liq}} R^2(0)}{\Delta P} \right)^{1/2}
\]
FIG. 3: The collapse time $t_{\text{clps}}$ for the metastable bubble in an overheated Fermi liquid as a function of relaxation time $\tau$ for different temperatures $T_0 = 4$ and 6 MeV shown near the curves at fixed asymmetry parameter $X_{\text{liq}} = 0.1$. The solid and dashed lines represent the calculations with non-Markovian and Markovian equations of motion respectively. The values of $\Delta T$ and $X_{\text{liq}}$ are the same as in Fig. 1.

with the yield $t_{\text{clps},R} = 6.1 \cdot 10^{-22}$ s which is in a good agreement with $t_{\text{clps,0}}$.

4. In conclusion we note that the collapse of the bubble in an overheated (undercompressed) Fermi liquid is strongly influenced by memory effects, if the relaxation time $\tau$ is large enough at $\omega_F \tau \gtrsim 1$. In this case, the collapse of the bubble is accompanied by characteristic shape oscillations of the bubble radius $R(t)$ (see Fig. 2) which depend on the memory kernel $K(t, t')$ and the relaxation time $\tau$. These oscillations appear due to the non-adiabatic elastic force induced by the memory integral. The non-adiabatic elastic force acts against the adiabatic force $-\partial \Phi/\partial R$ (see Eq. (6)) and hinders the collapse of the bubble. In contrast to the case of the Markovian motion, the delay in the collapse is caused here by both conservative elastic and friction forces.

We should like to point out that the study of the cavitation in a two-component nuclear
matter could represent a stimulating area of interest due to the fact that the vapor inside embryonic bubbles is strongly asymmetric system with $(N - Z)/(N + Z) \approx 1$, where $N$ and $Z$ are the number of neutrons and protons, respectively. The collapse of such extremely neutron rich system in which the energy is focused in a small region could show up new phenomena in the clusterization of nucleons.

The main goal of this paper was to show the existence of a new phenomena of shape oscillations of the collapsing bubbles in a hot nuclear Fermi liquid. This phenomena does not exist in classical (Newtonian) liquids. The shape oscillations of the collapsing bubble could induce a quite stimulating search of the accompanied characteristic $\gamma$-quanta emission. That provides, in principle, the possibility for the measurement of the temperature of the first kind phase transition through the measurement of the energy and the damping of the corresponding resonances in the $\gamma$-quanta spectrum.

Our consideration was restricted to a nuclear matter for maximum simplicity. Both effects of Coulomb forces and nuclear surface tension must be taken into account in the case of bubble dynamics in finite nuclei. The presence of Coulomb forces enhances the bubble formation and decreases the critical radius $R^*$. Note, however, that the minimum value of the radius $R^*$, for which our approach is applicable, is determined by the condition $a/R^* \ll 1$, where $a = 0.5 \div 1$ fm is the surface thickness of the bubble. The bubble dynamics in hot finite nucleus requires further investigations.

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