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Abstract

Early development of the instabilities in a dilute nuclear source is investigated using a finite temperature quantal RPA approach for different systems. The growth rates of the unstable collective modes are determined by solving a dispersion relation, which is obtained by parametrizing the transition density in terms of its multipole moments. Under typical conditions of a dilute finite system at moderate temperatures the dispersion relation exhibits an ultraviolet cut-off. As a result, only a finite number of multipole modes becomes unstable, and the number of the unstable collective modes increases with the size of the source. Calculations indicate that for an expanding source, unstable modes show a transition from surface to volume character.

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1 Introduction

The experimental observation of the abundant fragment production obtained in violent heavy ion collisions has generated many theoretical efforts for understanding the mechanisms responsible for such an explosion of the nuclear systems. It has been proposed that, because of the initial collisional shock, a large part of the nuclear matter phase diagram may be explored and new states can be accessed, hence making it possible to enter the unstable region of the phase diagram [1, 2, 3]. In such a context, fragment formation may take place through a rapid amplification of dynamical instabilities in the spinodal region. In order to deal with the dynamics of these large density fluctuations, stochastic semi-classical approaches of the *Boltzmann – Langevin* type, have been developed and applied to investigate the spinodal decomposition of nuclear systems [4, 5]. As far as the early development of instabilities is concerned, useful information can be gained, more easily, in the linear response framework of such approaches [6, 7]. Also, the response of the system to small initial perturbations can be studied within the Landau theory of Fermi liquid [8, 9]. It turns out that, due to the finite range of the nucleon-nucleon attraction, the small amplitude density inhomogeneities need to have a relative large spatial extension ($\approx 5-7\ fm$) in order to grow [2, 6]. The fact that the corresponding most unstable wave numbers ($k \approx 0.8 - 1\ fm$) are of the same order of magnitude as the Fermi momentum of the dilute systems suggests that quantal effects may have an important influence on the spinodal decomposition process, and should be included into the treatment for a quantitative description of the growth of instabilities. In a quantal RPA framework, it has been shown in [10] that in unstable nuclear matter, the most important modes shift towards longer wave lengths ($\lambda \approx 10\ fm$) due to quantal effects.

In this paper, we extend our previous quantal RPA treatment for the unstable nuclear matter, and investigate the early evolution of the unstable collective modes in the realistic case of the finite dilute systems, which may be formed during a nuclear collision. We calculate the growth rates of the unstable collective modes solving a quantal dispersion relation, and study the relative importance of the surface effect, the quantal effect and the finite range of the nucleon interaction in determining the most important unstable collective modes.

2 RPA Dispersion Relation

In the mean-field approximation, the single-particle density matrix ρ of the system is determined by the time-dependent Hartree-Fock (TDHF) equation,

$$i\hbar\frac{\partial\rho(t)}{\partial t} = [h[\rho],\ \rho(t)], \tag{1}$$

where $h[\rho] = \mathbf{p}^2/2m + U[\rho]$ denotes the mean-field Hamiltonian, and $U[\rho]$ is the density dependent self-consistent mean-field potential. The source is represented at the initial time $t = 0$ by a density matrix $\rho_0 = \rho(0)$ determined by the constraint Hartree-Fock equation $[h[\rho_0] - \lambda\hat{Q},\ \rho_0] = 0$, where $h[\rho_0]$ is the mean-field

Hamiltonian at the initial state, $\hat{Q}$ is a suitable constraining operator for preparing the system at low densities and λ is the associated Lagrange multiplier. The stability of a finite hot nuclear source against small amplitude density fluctuations has been studied on the basis of the TDHF eq.(1) in [11]. However, for this purpose it is more convenient to consider the density matrix $\hat{\rho}(t)$ in the "moving frame", $\hat{\rho}(t) = exp[\frac{i}{\hbar}\lambda t\hat{Q}]\ \rho(t)\ exp[-\frac{i}{\hbar}\lambda t\hat{Q}]$, and transform the TDHF equation into the moving frame,

$$i\hbar\frac{\partial\hat{\rho}(t)}{\partial t} = [\hat{h}(t) - \lambda\hat{Q},\ \hat{\rho}(t)], \tag{2}$$

where the mean-field Hamiltonian in the moving frame is given by,

$$\hat{h}(t) = exp[i\lambda t\hat{Q}]\ h(t)\ exp[-i\lambda t\hat{Q}]. \tag{3}$$

In order to investigate the early evolution of instabilities, we linearize this equation around $\hat{\rho}_0(t)$, $\hat{\rho}(t) = \hat{\rho}_0(t) + \delta\hat{\rho}(t)$, where $\hat{\rho}_0(t)$ is the solution of the TDHF eq.(2) with the initial condition ρ_0 determined by the constraint Hartree-Fock equation. The small fluctuation $\delta\hat{\rho}(t)$ is determined by the linearized TDHF equation in the moving frame,

$$i\hbar\frac{\partial\delta\hat{\rho}}{\partial t} = [\hat{h}_0(t) - \lambda\hat{Q},\ \delta\hat{\rho}] + [\delta\hat{U}(t),\ \hat{\rho}_0(t)] = \mathcal{M}(t)\cdot\delta\hat{\rho}(t), \tag{4}$$

where the mean-field Hamiltonian $\hat{h}_0(t)$ and the fluctuations of the mean-field potential $\delta\hat{U}(t)$ in the moving frame are defined in a manner similar to eq.(3). and $\mathcal{M}(t)$ denotes the instantaneous RPA matrix. The formal solution of this equation can be expressed as

$$\delta\hat{\rho}(t) = \mathcal{U}(t)\cdot\delta\rho(0), \tag{5}$$

where $\mathcal{U}(t) = \mathcal{T}(exp[-\frac{i}{\hbar}\int_0^t ds\mathcal{M}(s)])$ denotes the linearized evolution operator with $\mathcal{T}$ as the time ordering operator. The eigenvalues of the evolution operator $\mathcal{U}(t)$ determine the stability of the TDHF trajectories as a function of time. However the construction of $\mathcal{U}(t)$ is, in general, a very difficult task. Therefore, we consider the early evolution of the instabilities in the vicinity of the initial state ρ_0 and solve the RPA problem associated with $\mathcal{M}(0) = \mathcal{M}$. Introducing the eigenmodes $\delta\rho(\omega)$ associated with the eigenvalue $\hbar\omega$ and incorporating the representation $|i>$, which diagonalizes $h_0 - \lambda\hat{Q}$ and ρ_0, the RPA equation $\mathcal{M}\ \delta\rho(\omega) = \hbar\omega\ \delta\rho(\omega)$ for the collective modes becomes

$$(\hbar\omega - \epsilon_i + \epsilon_j)\ <i|\delta\rho(\omega)|j> = <i|\delta U(\omega)|j>\ (\rho_j - \rho_i), \tag{6}$$

where ρ_i and ϵ_i are the occupation number and the energy associated with the constraint Hartree-Fock state $|i>$, respectively. The temperature dependence enters into the calculations through the occupation number ρ_i, which is assumed to be given by the Fermi-Dirac function in terms of the single-particle energies ϵ_i.

The RPA eq.(6) can be solved using standard techniques [11, 12]. However, here we consider a simplified approach, and parametrize the transition density associated with an isoscalar collective mode in terms of its multiple moments as,

$$\delta\rho(\mathbf{r};\omega) =< \mathbf{r}|\delta\rho(\omega)|\mathbf{r} >= \alpha_L(\omega)\ f_L(kr)\ \bar{Y}_{LM}(\theta,\phi), \tag{7}$$

where $\alpha_L(\omega)$ is the amplitude, $f_L(kr)$ is the radial form factor associated with the collective mode, k is the radial wave number and $\bar{Y}_{LM}(\theta,\phi) = [Y_{LM}(\theta,\phi) + Y_{LM}{}^*(\theta,\phi)]/\sqrt{2(1+\delta_{M0})}$. We parametrize the radial form factor in terms of the Bessel function $j_L(kr)$ as $f_L(kr) = j_L(kr)\rho_0(r)$, where $\rho_0(r)$ is the equilibrium density of the source. By inverting this relation, the amplitude of the mode can be expressed according to,

$$\alpha_L(\omega) = K_L \int d^3r\ F_L(\mathbf{r})\delta\rho(\mathbf{r};\omega), \tag{8}$$

where $F_L(\mathbf{r}) = F_L(r)\bar{Y}_{LM}(\theta,\phi)$ is a function with a smooth r dependence and the normalization factor K_L is given by

$$\frac{1}{K_L} = \int d^3r\ F_L(\mathbf{r})\ f_L(kr)\ \bar{Y}_{LM}(\theta,\phi). \tag{9}$$

A dispersion relation for the frequencies of the collective modes can be deduced from the self-consistency condition that is obtained by inserting the solution of the RPA equation for $\delta\rho(\mathbf{r},\omega)$ into the right hand side of eq.(8). This gives

$$\frac{\alpha_L(\omega)}{K_L} = \sum_{i,j} \frac{\alpha_L(\omega)\ < i|\partial U/\partial\alpha_L|j >< j|F_L|i >}{\hbar\omega - \epsilon_i + \epsilon_j}\ (\rho_j - \rho_i), \tag{10}$$

where the transition field $\delta U_L(\omega)$ is written in terms of the collective amplitude $\alpha_L(\omega)$ as $\delta U_L(\omega) = (\partial U/\partial\alpha_L)\ \alpha_L(\omega)$. This dispersion relation is valid, in principle, for any choice of $F_L(r)$, provided that the parametrization (7) is a good approximation for the density fluctuations in a multipole mode. In fact, our investigations show that the solution of the dispersion relation is not very sensitive to the specific form of $F_L(r)$. Here, we take $F_L(\mathbf{r}) = \partial U/\partial\alpha_L$. This gives rise to a symmetric dispersion relation, which is equivalent to the RPA problem with a separable interaction of the form,

$$V(\mathbf{r},\mathbf{r}') = \frac{1}{2}\ \kappa_L\ F_L(\mathbf{r})F_L(\mathbf{r}'), \tag{11}$$

with the coupling constant κ_L given by the normalization factor K_L in eq.(9). Using the Wigner-Eckart theorem and the properties of the spherical harmonics functions, it is possible to sum over M_i, M_j and the dispersion relation becomes,

$$1 = K_L \sum_{n_i,L_i,n_j,L_j} 4\frac{\rho_j - \rho_i}{\hbar\omega - \epsilon_i + \epsilon_j} \frac{(2L_i+1)(2L_j+1)}{4\pi} \begin{pmatrix} L_j & L & L_i \\ 0 & 0 & 0 \end{pmatrix}^2 \tag{12}$$

$$| < i|\partial U/\partial\alpha_L|j >_R |^2.$$

Here $\begin{pmatrix} L_j & L & L_i \\ 0 & 0 & 0 \end{pmatrix}$ is a $3-j$ symbol and $< i|\partial U/\partial\alpha_L|j >_R$ denotes the radial matrix element of the potential.

The dispersion relation (12) allows to determine frequencies ω associated with the collective modes of the nuclear source. Since we consider a spherically symmetric source, collective frequencies depend on the multipole order L and the radial wave number k, but not on M. In the unstable region, the collective frequencies are imaginary and determine the growth rates $\tau_L(k)$ of the unstable collective modes, $\omega_L(k) = \pm i/\tau_L(k)$. In the calculations, we use a Skyrme-like parametrization for the effective mean-field potential,

$$U(\rho) = \frac{3}{4}\, t_0\, \rho + \frac{(\sigma+2)}{16}\, t_3\, \rho^{\sigma+1} + c\, \nabla^2\rho. \tag{13}$$

With the parameters $\sigma = 1$, $t_0 = 1000\ MeV\ fm^3$ and $t_3 = 1500\ MeV\ fm^6$ this force gives a saturation density of $0.16\ fm^{-3}$ and a compressibility of $350\ MeV$. The parameter $c = -126\ MeV\ fm^5$ is the same as in the corresponding term in the Skyrme-3 force. The radial matrix element in the dispersion relation (12) for a Skyrme-like effective mean-field potential can be expressed as

$$< i|\partial U/\partial\alpha_L|j >_R = \int r'^2 dr'\ (\partial U/\partial\alpha_L)_R\ h_{n_i,L_i}(r')\ h_{n_j,L_j}(r'), \tag{14}$$

where $h_{n_i,L_i}(r)$ denotes the radial wave function and

$$(\partial U/\partial\alpha_L)_R = [\frac{3}{4}\, t_0 + \frac{(\sigma+1)(\sigma+2)}{16}\, t_3\, \rho^{\sigma} + c\, (\frac{L(L+1)}{r^2} + \frac{1}{r^2}\frac{d}{dr}r^2\frac{d}{dr})\ ]\ f_L(kr). \tag{15}$$

3 Instabilities in Finite Systems

In order to solve the dispersion relation (12), first we need to determine the single-particle representation of the constraint Hartree-Fock problem (CHF). However, it is not very easy to guess a suitable form of the constraining operator $\hat{Q}(r)$, which produces the single-particle representation for a wide range of densities in the unstable region. For example, an operator of the form

$$\hat{Q}(r) = r^2\, e^{-r/r_0}, \tag{16}$$

where r_0 is a cut-off distance, provides a resonable constraint for preparing the nuclear source at low densities. However, in the spinodal region, as soon as the monopole mode $L = 0$ is unstable, the CHF calculations can not be carried out. For this reason, we follow a more schematic approach and solve the dispersion relation by employing the harmonic oscillator wave functions and the wave functions of a Wood-Saxon-like potential, instead of the CHF wave functions. The harmonic oscillator representation implicitly corresponds to a constraint according to,

$$U[\rho] - \lambda\hat{Q}(r) = \frac{1}{2}Cr^2 \tag{17}$$

where C is the stiffness parameter of the oscillator determined by the root-mean-square-radius of the source. The representation of the Wood-Saxon-like potential is generated by solving the eigenfunctions of the Skyrme potential $U(\rho_0)$ taken as a function of the source density ρ_0. The source density profile is taken as a Fermi shape $\rho_0(r) = \bar{\rho}(0)/[1+exp(r-R)/a]$, where $\bar{\rho}(0)$, R and a are the central density, sharp radius and the surface thickness of the source, respectively. In particular, this representation provides a useful basis to investigate the influence of the source density profile on the properties of the unstable modes. In order to obtain accurate solutions of the dispersion relation, a sufficiently large number of orbitals should be incorporated into the calculations. Here, we present calculations carried out for sources containing $A = 40$ and $A = 140$ nucleons by including 100 and 120 orbitals, respectively.

In the top panel of figure 1, minimum values of the quantity $\omega^2/|\omega|$ for modes with multipolarity $L = 0, 2, 3$ are plotted for a source containing $A = 40$ nucleons as a function of the root-mean-square-radius $< r^2 >^{1/2}$ of the system at zero temperature. These results are obtained by solving the dispersion relation employing the harmonic oscillator representation (left panel) and the CHF representation (right panel) with a constraint given by eq.(16). When the minimum value of the quantity $\omega^2/|\omega|$ is negative, the corresponding mode becomes unstable. In this manner the CHF calculations provide a good basis for understanding transition between the stable and the unstable regions . As seen from the upper right part of figure 1, for increasing root-mean-square-radius the collective modes become softer and around $< r^2 >^{1/2} \approx 3.8\ fm$ the octupole mode becomes unstable. As indicated above, the CHF calculations are not reliable for $< r^2 >^{1/2} \geq 4.1\ fm$ when the system becomes too dilute. However, as seen from the upper left part of figure 1, the results obtained with the harmonic oscillator representation compares rather well with those of the CHF calculations. In both calculations, the system begins to show up instabilities when the root-mean-square-radius of the source is around 3.7-3.8 fm. We notice that the instability growth rates obtained in both calculations for L =3 mode compare rather well. Moreover, the calculated frequencies of the stable oscillations for L =0,2,3 modes are also quite similar in both calculations. In the bottom panel of figure 1, the density profiles of the source corresponding to two different root-mean-square-radii are shown in both calculations. Figure 2 illustrates the effect of the Coulomb force on the instabilities of the most important unstable modes for sources with $A = 40$ and $A = 140$ nucleons at zero temperature. The Coulomb force is included into the calculations by replacing $\partial U/\partial \alpha_L$ with $\partial U/\partial \alpha_L + (1/2)\partial U^C/\partial \alpha_L$ in the dispersion relation (12). Here, the fluctuating part of the mean-field due to the Coulomb force is calculated according to [13],

$$\delta U^C(\mathbf{r}) = \int \frac{e^2}{|\mathbf{r}-\mathbf{r}'|}\delta\rho_p(\mathbf{r}')\, d^3r' \; - \frac{1}{3}\left(\frac{3}{\pi}\right)^{1/3} e^2\, \rho_p^{-2/3}(\mathbf{r})\, \delta\rho_p(\mathbf{r}), \qquad (18)$$

with $\delta\rho_p = \delta\rho/2$, where the second term is the Slater approximation for the exchange part of the interaction. In figure 3, the minimum values of $\omega^2/|\omega|$ for quadrupole and octupole modes are plotted for the same systems as a function of the root-mean-square-radius of the system at temperatures $T = 0\ MeV$ and $T = 3\ MeV$. These calculations in figure 2 and figure 3 are performed in the harmonic oscillator representation. As seen from the figures, the Coulomb force has a minor effect in the light system, and in the case of the heavier system, the degree of the instability is slightly decreased by the Coulomb force. On the other hand, a hot source is more stable than the source at $T = 0\ MeV$, and for increasing temperatures the source becomes unstable at more dilute configurations than those of at zero temperature.

Figure 4 shows the radial wave numbers associated with the quadrupole and octupole modes as a function of the root-mean-square-radius of the density distribution in a source with $A = 40$ nucleons (left part) and $A = 140$ nucleons (right part) at zero temperature, calculated with the harmonic oscillator representation. These wave numbers correspond to the lowest mode in the stable region and the most unstable mode in the unstable region. In the figure, the crossover from stable to unstable regions is indicated by vertical lines. In the case of the quadrupole mode, the crossover occur at $< r^2 >^{1/2} \approx 4.2\ fm$, whereas the octupole mode becomes unstable already at $< r^2 >^{1/2} \approx 3.7\ fm$. In both cases, modes exhibit a rather rapid transition from surface character with small values of k to volume character with large values of k at around $< r^2 >^{1/2} \approx 4.0 - 4.5\ fm$. It is also seen that the transitions from stable to unstable regions and from surface to volume character occur at different stages.

In the bottom panel of figure 5, the dispersion relations for different multipolarities are plotted as a function of the radial wave number k for sources containing $A = 40$ nucleons (left part) and $A = 140$ nucleons (right part) at a temperature $T = 3\ MeV$. The results obtained in the harmonic oscillator and the Wood-Saxon representations are indicated by solid-lines and dashed-lines, respectively. The density profiles associated with the sources are displayed in the top panel of the figure. As seen from the figure, except for the lowest unstable mode, dispersion relation is not very sensitive to the representation employed. We note also that, the growth rates for large values of the radial wave number are suppressed due to the quantal and surface effects, indicating that the system does not exhibit more than one radial oscillation in the unstable modes. In the case of $A = 40$, the calculations done with the harmonic oscillator representation show that the system is unstable against quadrupole and octupole deformations . On the other hand, in the calculations with the Wood-Saxon representation the octupole is the dominant unstable mode. This is consistent with the CHF calculations presented in figure 1, in which the lowest mode that becomes unstable when the system is diluted is the octupole mode. In the system with $A = 140$ nucleons, several high order multiple modes up to $L = 5$ become unstable. Due to the quantal and surface effects, the multipoles with L larger than 5 are strongly suppressed. The maximum value of the frequencies $|\omega_L(k)|$ is nearly equal for all multipoles in the region $L = 2 - 5$, indicating that these modes can be excited, apart from the statistical weight $2L + 1$, with nearly equal probability [14]. These results are in agreement with recent calculations based on a fluid dynamic approach to spinodal instabilities [15]. It is interesting to note that the maximum of the growth rate for a typical multipole mode in a finite source is comparable to the one obtained in nuclear matter. In fact, we perform a calculation in a periodic box by solving eq.(12) and determine the growth rates of the unstable modes as a function of the wave number. We find that the maximum growth rate at a given density is close to the growth rate of a typical mode in a finite system at the same central density.

In figure 6, the maximum value of the frequency $|\omega_L|$ obtained in the calculations using the harmonic oscillator representation is plotted as a function of L for $A = 40$ (left part) and for $A = 140$ (right part) at temperatures $T = 0, 3, 5\ MeV$. These calculations correspond to a source with the root-mean-square radius taken as $4.54 fm$ for the small system and $6.21 fm$ for the large system. It is seen that the instability in both systems decreases for increasing temperature of the source as expected, and the octupole mode appears, once again, as the most robust one. Also, in the large system, the dispersion relation has a cut at a lower multipolarity for increasing temperature.

4 Conclusions

In order to investigate the early development of instabilities in a dilute nuclear source, we carry out finite temperature quantal RPA calculations for systems with $A = 40$ and $A = 140$ nucleons. A parametrization of the transition density in terms of its multipole moments leads to a simple dispersion relation for the growth rates of the unstable collective modes . We determine the growth rates as a function of the radial wave number from the dispersion relation employing a suitable single-particle representation. Under typical conditions, when the dilute system with $A = 140$ nucleons has an average density $\rho = 0.05 fm^{-3}$ and a temperature range $T = 3 - 5 MeV$ the collective modes up to $L = 5 - 6$ become unstable. Furthermore, as the source expands to lower densities, the unstable modes exhibits a transition from surface to volume character. The maximum growth rates of these unstable modes are nearly the same around 30 fm/c, indicating that the system may develop into different fragmentation channels with nearly equal probability. The results presented here are consistent with recent calculations of spinodal instabilities in finite nuclear systems based on a fluid dynamic approach.

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Figure 1:

Top panel: minimum values of $\omega^2/|\omega|$ for modes with multipolarity $L = 0, 2, 3$ for $A = 40$ as a function of the root-mean-square-radius $< r^2 >^{1/2}$ at zero temperature calculated in the harmonic oscillator representation (left part) and the constraint Hartree-Fock representation (right part). Bottom panel: density profile of the source corresponding to different root-mean-square-radii in the harmonic oscillator representation (left part) and the constraint Hartree-Fock representation (right part).

Figure 2:

Minimum values of $\omega^2/|\omega|$ for different multipole modes for $A = 40$ (left part) and $A = 140$ (right part) as a function of the root-mean-square-radius with the Coulomb force (dashed lines) and without the Coulomb force (solid lines) at zero temperature, calculated in the harmonic oscillator representation.

Figure 3:

Minimum values of $\omega^2/|\omega|$ for quadrupole and octupole modes for $A = 40$ (left part) and $A = 140$ (right part) as a function of the root-mean-square-radius at temperatures $T = 0\ MeV$ (dashed lines) and $T = 3\ MeV$ (solid lines), calculated in the harmonic oscillator representation.

Figure 4:

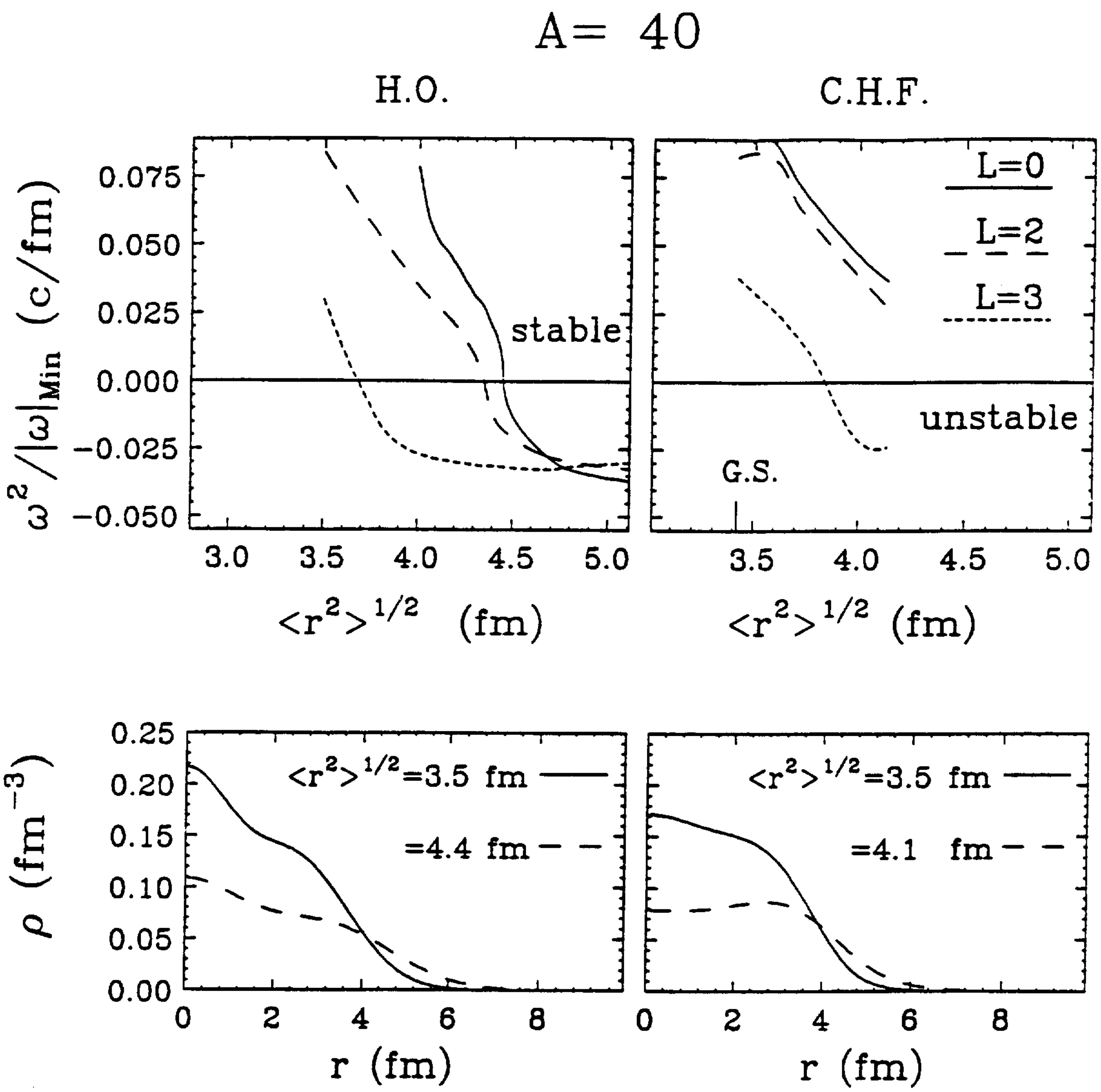
The radial wave numbers associated with the quadrupole mode (top panel) and the octupole mode (bottom panel) as a function of the root-mean-square-radius in a source with $A = 40$ at zero temperature. The vertical lines indicate the crossover from stable to unstable regions.

Figure 5:

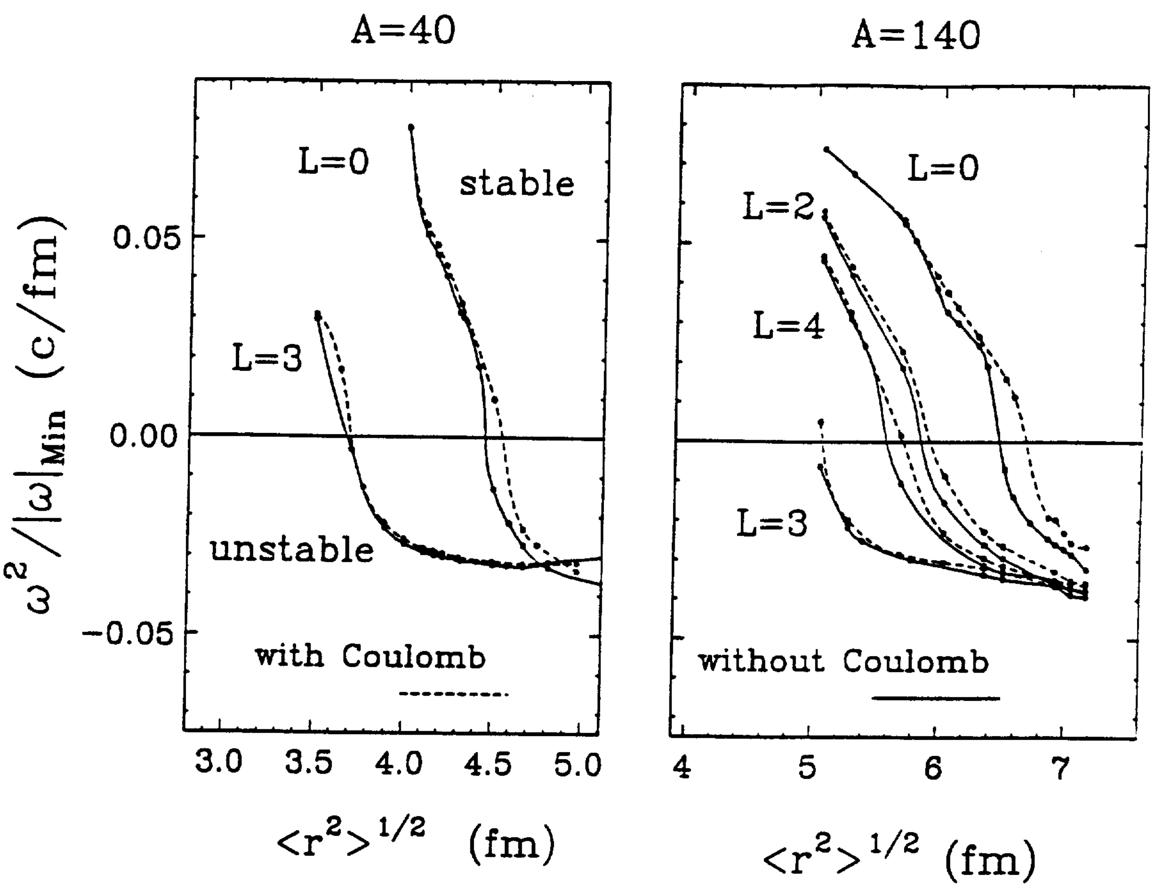
Top pannel: dispersion relations calculated in the harmonic oscillator representation (solid lines) and the Wood-Saxon representation (dashed lines) for different multipolarities plotted as a function of the radial wave number k for $A = 40$ nucleons (left part) and $A = 140$ (right part) at a temperature $T = 3$ MeV. Bottom panel: density profile of sources in the harmonic oscillator representation (solid lines) and the Wood-Saxon representation (dashed lines).

Figure 6:

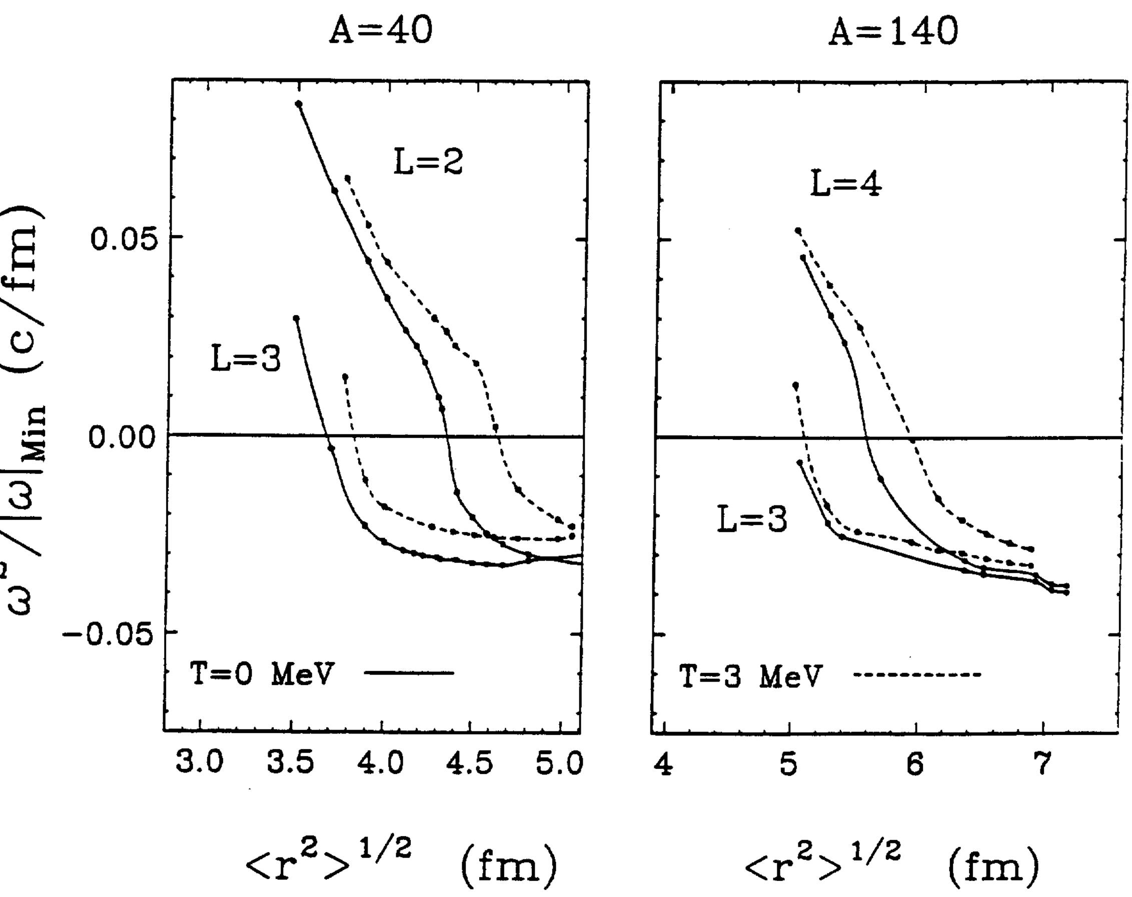
The maximum value of the frequency $|\omega_L|$ obtained in the harmonic oscillator representation as a function of the multipolarity L for $A = 40$ (left part) and for $A = 140$ (right part) at temperatures $T = 0, 3, 5\ MeV$.

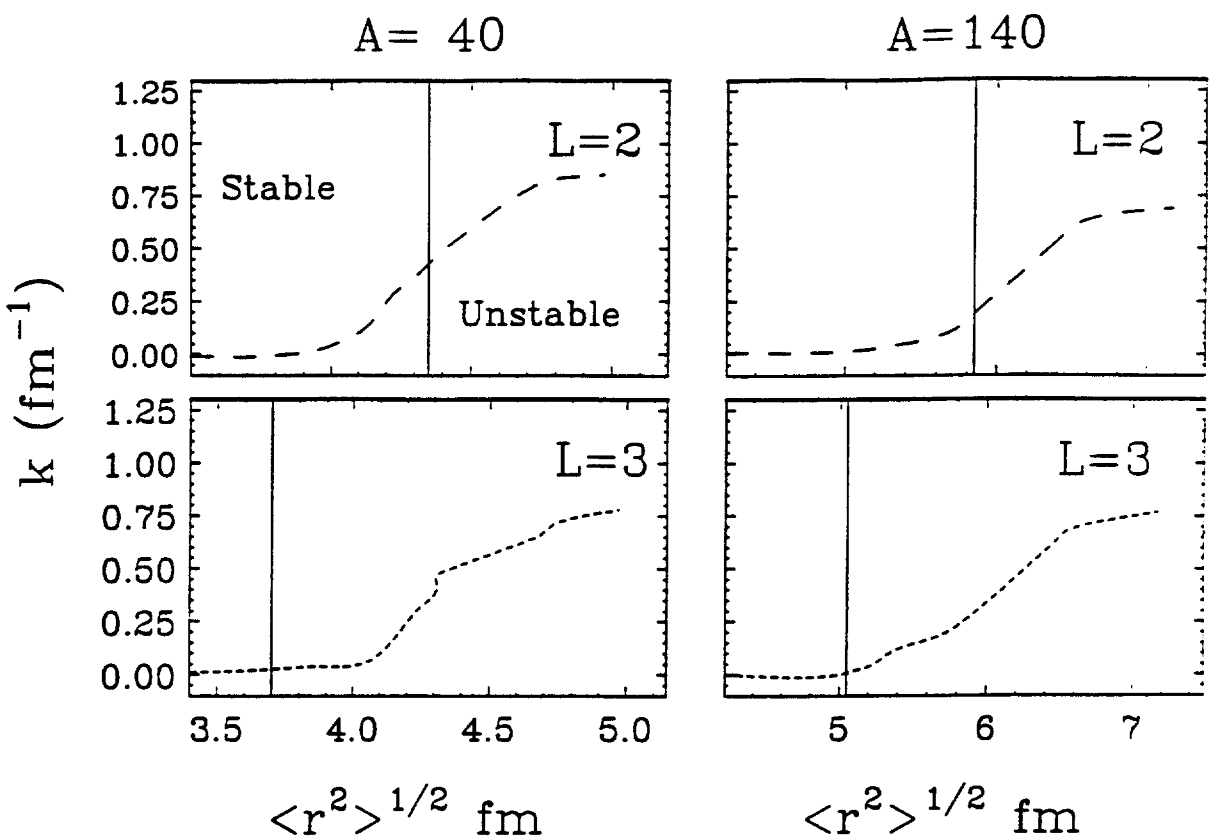


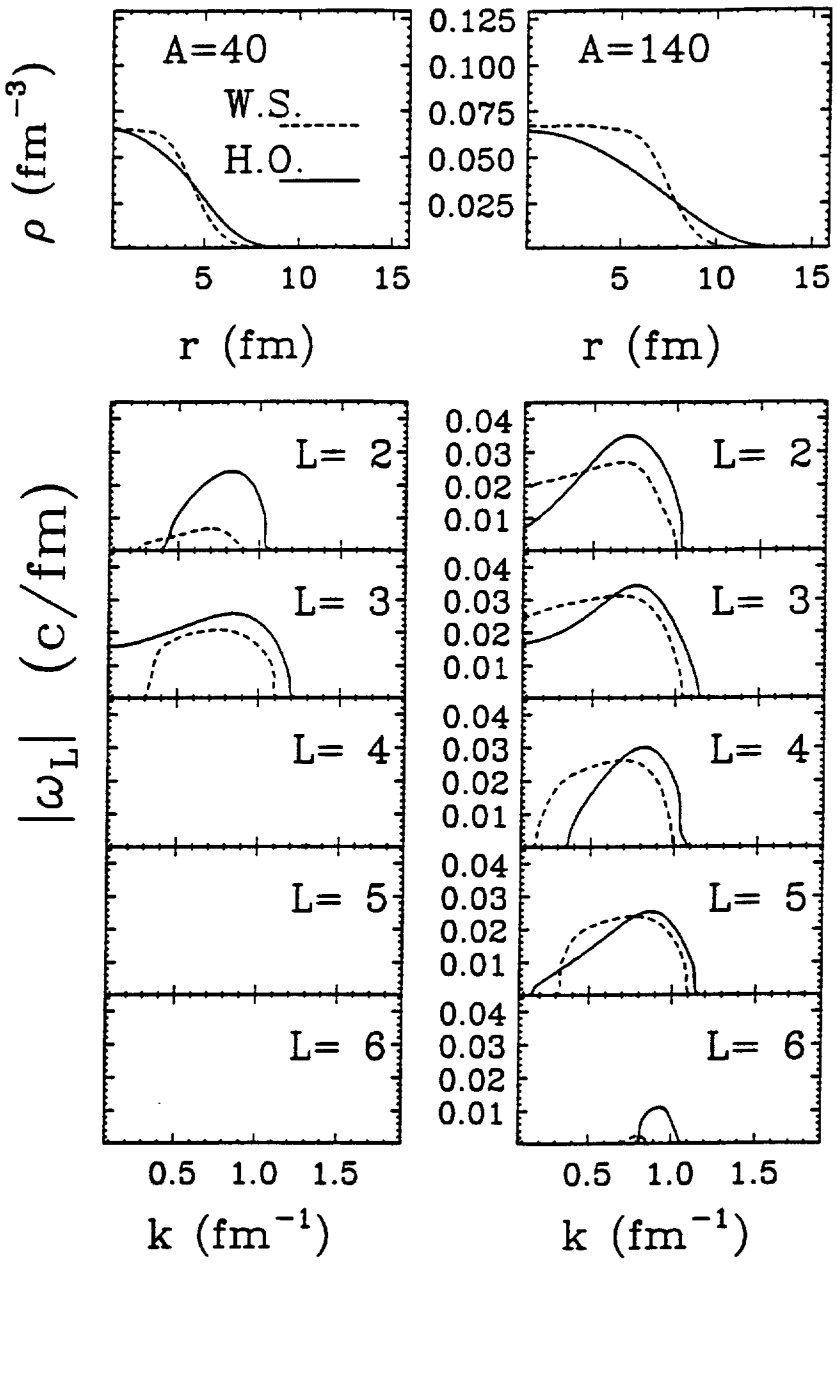
A= 40
H.O.
C.H.F.
L=0
L=2
L=3
stable
unstable
G.S.
$\omega^2/|\omega|_{Min}$ (c/fm)
0.075
0.050
0.025
0.000
−0.025
−0.050
3.0
3.5
4.0
4.5
5.0
$<r^2>^{1/2}$ (fm)
$<r^2>^{1/2}$=3.5 fm
=4.4 fm
=4.1 fm
ρ (fm^{-3})
0.25
0.20
0.15
0.10
0.05
0.00
0
2
4
6
8
r (fm)

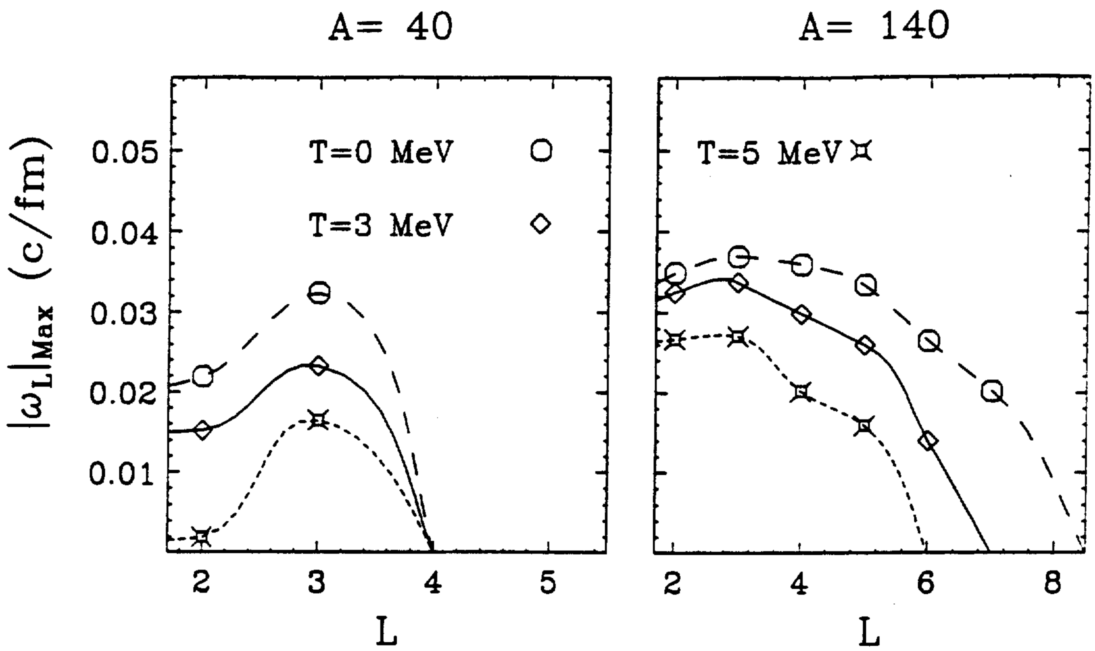


A=40
A=140
L=0
stable
L=3
unstable
with Coulomb
L=0
L=2
L=4
L=3
without Coulomb
0.05
0.00
−0.05
$\omega^2/|\omega|_{Min}$ (c/fm)
3.0
3.5
4.0
4.5
5.0
4
5
6
7
$<r^2>^{1/2}$ (fm)
$<r^2>^{1/2}$ (fm)


A=40
A=140
L=2
L=3
L=4
L=3
0.05
0.00
−0.05
$\omega^2/|\omega|_{Min}$ (c/fm)
T=0 MeV
T=3 MeV
3.0
3.5
4.0
4.5
5.0
4
5
6
7
$<r^2>^{1/2}$ (fm)
$<r^2>^{1/2}$ (fm)


A= 40
A=140
k (fm^{-1})
Stable
Unstable
L=2
L=2
L=3
L=3
1.25
1.00
0.75
0.50
0.25
0.00
3.5
4.0
4.5
5.0
5
6
7
<r^2>$^{1/2}$ fm
<r^2>$^{1/2}$ fm


A=40
W.S.
H.O.
A=140
ρ (fm⁻³)
0.125
0.100
0.075
0.050
0.025
5
10
15
r (fm)
L= 2
L= 3
L= 4
L= 5
L= 6
0.04
0.03
0.02
0.01
|ω_L| (c/fm)
0.5
1.0
1.5
k (fm⁻¹)



A= 40
A= 140
$|\omega_L|_{Max}$ (c/fm)
T=0 MeV
T=3 MeV
T=5 MeV
0.05
0.04
0.03
0.02
0.01
2
3
4
5
L
2
4
6
8
L