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Introduction

It is with great pleasure that I introduce this volume to the scientific community interested in microgravity research. It contains selected papers from our most recent IXth European Symposium on "Gravity-dependent Phenomena in Physical Sciences", which was held in Berlin in May 1995. It was organized by ESA together with the German Space Agency DARA and the Technical University Berlin. The international audience present at this conference as well as the quality of the papers and of the discussions is a clear demonstration to me that our investment in this field of research and applications is bearing fruit and that we have reached a level of maturity which is well up to international standards.

I am particularly satisfied with this development in view of the far-reaching decisions taken at the ESA Council meeting at Ministerial Level in Toulouse a few weeks ago, which will shape the future of ESA's programmes for decades to come. One of the major decisions was on the European participation in the International Space Station Programme. Recognizing the significant scientific, technological and political benefits that will accrue to Europe from our participation in this programme, the Ministers decided the financing and the development of the Columbus Orbiting Facility and the Automated Transfer Vehicle to be launched by Ariane-5 to service the Station. They also approved definition studies for a Crew Transport Vehicle and preparation activities for Station utilization including a programme for the development of Microgravity research facilities for Columbus. The Station exploration programme, which will run from 2002 to 2013 was also very clearly defined.

Microgravity research and applications will therefore have ample opportunities in the two decades to come and I am fully confident that Microgravity together with Earth Observation, Space Sciences and Technology will seize these opportunities and endeavour into a multidisciplinary and international utilization of this unique space infrastructure.

Paris, November 1995

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Theoretical Models for Boiling at Microgravity

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Abstract

In this work initial results are presented on the theoretical modelling of the following aspects related to microgravity boiling: the rate of phase change at moving interfaces, the thermocapillary effect at the phase separating boundary, the influence of thermal effects on the Rayleigh – Taylor instability of a heated static vapour layer below its liquid and the formation of the liquid macrolayer beneath vapour bubbles.

1. Introduction

Despite many years of experimental investigation the predictability of microgravity boiling is yet impossible [1]. The classical formulas for peak heat flux [2, 3, 4] for example, can no longer define its actual values at reduced gravity levels. Indeed, the existing theoretical models have been specifically developed for ground-based applications. The critical wavelength of the Rayleigh – Taylor instability [5] is used [3, 4] for the estimation of the distance between vapour masses on a plane heater in the stage of developed nucleate boiling. The critical wavelength has a small value at earth conditions and becomes very large as the gravity level decreases (for the water/air system under gravity $10^{-4}g_0$, $g_0 = 981 cm/s^2$) its value is equal to 160 cm). When the horizontal dimensions of a system are less then the value of the critical wavelength, the Rayleigh – Taylor instability does not occur since all perturbations are stabilized by surface tension. Because of the small heater areas used in the microgravity experiments performed up to now the existing theoretical models cannot be applied.

The focus of this work is related to interactions between the mechanisms of the thermocapillary and the Rayleigh – Taylor instability at the liquid-vapour interface in the case of low Bond numbers. The nonequilibrium rate of phase change is described using the irreversible thermodynamics. The influence of the nonequilibrium effect on the thermocapillary one is described in terms of the phenomenological mass transport coefficient. It is proven that at high heat flux in a very thin vapour layer the Rayleigh – Taylor instability is replaced by

the thermocapillary one if the mass transport coefficient is sufficiently small. A possible mechanism of the liquid macrolayer formation is discussed which is due to the thermocapillary instability of a residual vapour layer remaining on a heater after bubble break off.

2. Rate of phase change

To describe theoretically boiling in microgravity it is necessary to know precisely the thermal conditions at the interfaces. In studies on condensation and evaporation the rate of phase change is generally defined by the famous Hertz – Knudsen equation. For a flat interface, along which the liquid and the vapour temperatures are assumed to be equal, the Hertz – Knudsen equation [6] is

$$J = \beta(M/2\pi RT)^{1/2}[p_{\rm s}(T) - p_0(T)] \tag{1}$$

where J is the nonequilibrium mass flux across the interface, M is the molecular weight of the vapour, $p_{\rm s}(T)$ is the saturation pressure at surface temperature T, $p_0(T)$ is the vapour pressure just beyond the interface, R is the universal gas constant and β is the constant accommodation coefficient. According to the theoretical considerations of the kinetic theory the value of β should be close to unity. However, reported experimental values are of order $10^{-3} - 10^{-1}$ and there also exist discrepancies among the experimental data of various investigations.

One can derive an equation for the kinetics of phase transition in the framework of irreversible thermodynamics. Assuming continuity of the temperature along the interface $T_{\rm v} = T_\ell \quad (= T)$ Onsager's law suggests [7] that

$$J = K[\mu_\ell(p_\ell, T) - \mu_{\rm v}(p_{\rm v}, T)] \ . \tag{2}$$

Here K is the positive phenomenological coefficient, $\rho_{\rm v}$ and ρ_ℓ are the densities of vapour and liquid respectively. J is the rate of phase change. By definition $J = \mathbf{J} \cdot \mathbf{n}$, where $\mathbf{J} = \rho_\ell(\mathbf{v}_\ell - \mathbf{v}_\Sigma) = \rho_{\rm v}(\mathbf{v}_{\rm v} - \mathbf{v}_\Sigma)$ is the mass flux vector at the interface, $\mathbf{v}_\ell$ and $\mathbf{v}_{\rm v}$ are the velocities of the liquid and vapour at the interface, $\mathbf{v}_\Sigma$ is the interface velocity and $\mathbf{n}$ is the surface normal unit vector directed into the vapour phase. Equation (2) together with the definition of J show that the mass will leave the phase with the highest chemical potential toward that with the lowest one. Let us remark that if the temperature jump across the liquid-vapour interface is taken into account, then other transport coefficients (see [8]) appear in the phenomenological equation for the mass and energy fluxes because there exists then a new thermodynamic force $\delta T = T_\ell - T_{\rm v}$. In accordance with the requirements of irreversible thermodynamics the phenomenological coefficient K depends on the surface temperature T and the vapour pressure $p_{\rm v}$. Equation (2) can be interpreted as an approximation to the general mass flux equation where the thermodynamic force δT is neglected.

In (2) the liquid and vapour pressures p_ℓ and $p_{\rm v}$ and the temperature T vary along the moving interface. Introduce now a function $p_{\rm s}(T)$ which relates the value of T to the corresponding saturation vapour pressure for a flat isothermal interface at a constant temperature T . Assume further that for the liquid

and vapour phases the equations of state are respectively $\rho_\ell = \rho_\ell^0(p, T)$, $\rho_v = \rho_v^0(p, T)$ and that everywhere along the liquid-vapour interface one has the inequalities $|p_\ell - p_s(T)|/p_s(T) \ll 1$, $|p_v - p_s(T)|/p_s(T) \ll 1$. One then can linearize the function K and the chemical potentials in the vicinity of the point $(p_s(T), T)$ using a Taylor expansion. With the assumption of local thermodynamic equilibrium inside each phase near the interface one can use the classical thermodynamical relation for the pressure derivative of the chemical potential [7]. The one obtains a linearized form of (2)

$$J = K(T)\left[\frac{p_s(T) - p_v}{\rho_v(T)} - \frac{p_s(T) - p_\ell}{\rho_\ell(T)}\right], \tag{3}$$

where $\rho_v(T)$, $\rho_\ell(T)$ and $K(T)$ stand for $\rho_v^0(p_s(T), T)$, $\rho_\ell^0(p_s(T), T)$ and $K(p_s(T), T)$ respectively. Far from the critical boiling point the vapour density is much lower than the density of liquid; assuming also $p_v \sim p_\ell$ and neglecting the last term in (3), one obtains $J = K(T)[p_s(T) - p_v]/\rho_v(T)$. This equation presents a formal analogy with the Hertz-Knudsen equation (1) since they are both proportional to the difference $p_s(T) - p_v$. But it has been deduced in a different way, without the ideal gas approximation for the vapour phase. In this equation one can use the correct state equation, experimentally evaluated for the vapour; assuming the interface is flat and isothermal, one can compare this equation with (1) and obtains the expression of the phenomenological coefficient K through the accommodation coefficient β :

$$K = \beta \rho_v(T)(M/2\pi RT)^{1/2}. \tag{4}$$

This expression is rough because it is obtained with the help of the linearization of the original (2). The nonlinear terms, neglected in (3), remain yet to be estimated. Probably, the investigation of the nonlinear approximations to (2) and its extension to the case where the temperature jump δT is taken into account will help to understand why so large discrepancies exist among the reported experimental data on the accommodation coefficient values. Note that in the nonlinear approximations the dependence of the phenomenological coefficient on the vapour pressure cannot be neglected.

The last term in (3) can be important for strongly curved interfaces. For this case, using the interfacial balance of normal momentum, we obtained from (3) the following approximate formula (H - is the mean curvature of moving interface) :

$$J = \frac{\rho_\ell(T) - \rho_v(T)}{\rho_\ell(T)\rho_v(T)}K(T)\left[p_s(T) - p_v - \sigma\frac{\rho_v(T)}{\rho_\ell(T) - \rho_v(T)}H\right]. \tag{5}$$

Equation (5) takes into account the influence of the surface tension on the rate of phase change.

3. The degree of nonequilibrium number

The rate of phase change depends on the interfacial temperature (2). Therefore, the thermocapillary force and the surface gradient of the nonequilibrium mass flux are interrelated. The inverse coefficient K^{-1} measures the deviation from the equilibrium state of the interface. Indeed, in the limiting case $K^{-1} \to 0$ it follows from (2) that the chemical potentials are equal: the moving interface is in a state of quasi-equilibrium. The influence of nonequilibrium mass transfer on the interfacial balances of momentum and energy is increased with increasing values of K^{-1}. In the case where $K^{-1} \to \infty$ the equality $J = 0$ obtained from (2) is the non-leak condition at the interface. Consequently, the thermocapillary problem with no phase change corresponds to this latter case.

The behavior of a liquid-vapour system can be quite different in the opposite cases "no phase change" and "quasi-equilibrium". We shall demonstrate this by examining the linear stability of a motionless liquid-vapour system (Fig. 1). A horizontal vapour layer lies between a wall being laterally infinite and which is heated at a constant temperature T_w and a liquid phase which extends infinitely. There is no mass transfer across the interface in the basic state, but it may occur in the perturbed state. We consider the heat conduction and dynamics of both phases, assuming them to be incompressible viscous fluids. Equation (3) is linearized at the saturation temperature T_0 of the unperturbed interface and is used as an interfacial boundary condition. This physical set-up allows one to investigate the interplay between the thermocapillary, the nonequilibrium phase exchange and the Rayleigh - Taylor mechanisms.

Our particular problem shows that the influence of the thermocapillary effect on the Rayleigh - Taylor instability depends essentially on the degree of nonequilibrium of the perturbed interface. The variation of the nonequilibrium mass flux along the interface gives a contribution to the total thermocapillary force characterized by the following dimensionless parameter,

$$s_n = \sigma_T T_0 / (K_0 L \nu_v).$$

Here $\sigma_T > 0$ is the temperature coefficient of surface tension, ν_v is the kinematic viscosity of the vapour, L is the latent heat of vaporization and $K_0 = K(T_0)$. Since s_n is inversely proportional to the phenomenological coefficient, it is a measure of the degree of interface nonequilibrium. Thermocapillary instability may exist only for high values of this parameter.

4. Results of the stability analysis

a) When s_n is infinite, the phase change effect is completely ignored. The problem is reduced to the study of the coupling between the pure Marangoni and the Rayleigh - Taylor instability mechanisms. The relative contribution of the destabilizing gravity mechanism to the total surface deformation effect is char-. acterized by the dimensionless group

$$G = k_* N_g / (24 N_{cr}) = (\rho_\ell - \rho_v) g d_v^3 / (24 \rho_v \nu_v k_\ell)$$
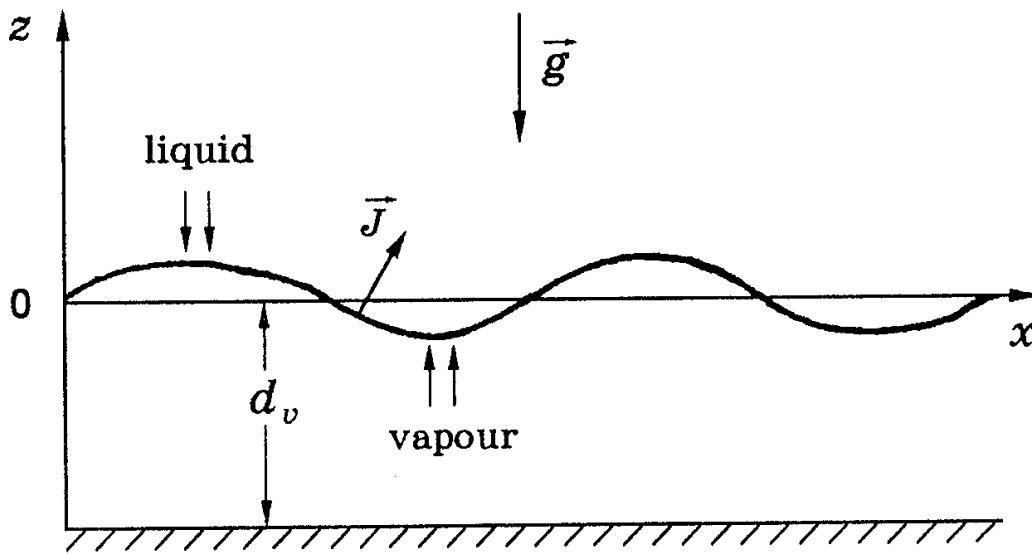
Fig. 1. Physical configuration of the liquid – vapour system

where $k_* = k_v/k_\ell$ is the thermal diffusivity ratio, d_v is the depth of the vapor layer. The parameters $N_g = (\rho_\ell - \rho_v)gd_v^2/\sigma_0$, $N_{cr} = \rho_v\nu_v k_v/(\sigma_0 d_v)$ are the Bond number and the crispation group. σ_0 is the value of the surface tension at $T = T_0$. If $G \ll 1$ the Rayleigh – Taylor instability is replaced by the thermocapillary surface-wave instability, in the case of a common liquid – vapour systems sufficiently heated from below,. This is demonstrated in Fig. 2. The marginal non–oscillatory instability curves are presented for the water system at $T_0 = 373K$. Note that in this "no phase–change" case the Marangoni number can be used to characterize the instability. However, the result is presented with respect to the number $Q = \lambda_v(T_w - T_0)/(\rho_v k_v L)$ -modified Jakob number, because it is more convenient for comparison with the "quasi-equilibrium" case, λ_v is the vapour thermal conductivity. For $G = 2.44 \times 10^{-6}$ the number Q is plotted versus the dimensionless wave number ω (the length scale is d_v). The curve 1 intersects the axis $Q = 0$ at $\omega = \sqrt{N_g} = 1.98 \times 10^{-5}$, the crossing point is not seen because it is very close to the coordinate origin. In the case $G \ll 1$ for large positive Q the condition of instability is given by the inequality $\omega < \omega_a$, where $\omega_a = \sqrt{24N_{cr}/k_*}$ is the position of the asymptote. The ratio between the actual critical wavelength $(2\pi d_v/\omega_a)$ and the critical wavelength $(2\pi d_v/\sqrt{N_g})$ of the pure Rayleigh – Taylor instability (which pertains to isothermal systems) is equal to the square root of the small parameter G. Note that the condition $G \ll 1$ is equivalent to the following inequality for the vapor depth $d_v \ll d^* = (24\rho_v\nu_v k_\ell)^{1/3}/([\rho_\ell - \rho_v)g]^{1/3}$. For the water system at earth gravity acceleration, $d^* = 1,7 \times 10^{-3}$ cm.

b) An essential difference of the "quasi-equilibrium" case ($s_n = 0$) with the previous one is that the thermocapillary instability is now absent and that the Marangoni number cannot be introduced. Here the thermocapillary number

$$s_e = \sigma_T T_0/(\rho_v L d_v)$$

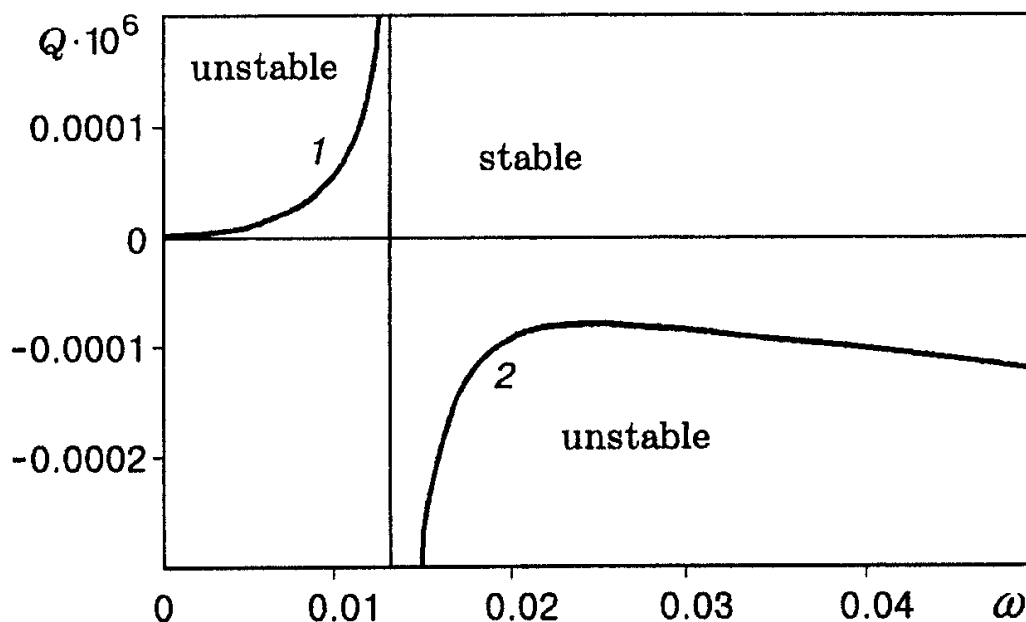characterizes the thermocapillary effect. For $s_e \ll 1$ the influence of the ther-

Fig. 2. Marginal stability curves for the water system at $T_0 = 373K$ in the case of no phase-change, $g = 981 \times 10^{-4} cm/s^2$, $d_v = 5 \times 10^{-4}$ cm

mocapillary effect on the neutral stability condition is shown to be negligible. The condition $s_e \ll 1$ is equivalent to the following inequality $d_v \gg d_* = \sigma_T T_0/\rho_v L$. The value of d_* is very small. Indeed for the water system at $373K$, $d_* = 5.5 \times 10^{-6}$ cm.

Our analytical investigations show that for the system heated from below ($Q > 0$) the phase change mechanism stabilizes the long-wave perturbations. In contrast with the previous case, the positive number Q now characterizes the stabilizing influence of the change of phase mechanism. With small G the Rayleigh – Taylor instability is prevented for sufficiently large positive Q. Such a situation is shown in Fig. 3 for the case when $G = 0.195$.

There is one single marginal stability curve which intersects the axis $Q = 0$ at $\omega = \sqrt{N_g}$ and at $\omega = 0$. The maximal value Q_m on this curve is equal to 7.2×10^{-9}. It is reached for $\omega = 0.00078$. For $Q > Q_m$ the motionless state of system is stable with respect to all monotonical perturbations.

c) Consider the case where the value of s_n is finite and positive. It is noteworthy that the thermocapillary effect is now characterized by two dimensionless parameters: s_n and s_e. Assume $G \ll 1$, $Q > 0$. Then owing to the disappearance of gravity waves, instability can be generated only by the thermocapillary mechanism. Instead of Q, the Marangoni number may be formed from dimensionless groups by the thermocapillary number s_e. However, the possible existence of thermocapillary instability depends on the value of the dimensionless group

$$\alpha = k_*/(N_{cr}s_n) = \sigma_0 K_0 L d_v/(\rho_v k_\ell \sigma_T T_0).$$

For high values of Q the instability criteria is formulated only in terms of the number α as follows. For $s_e \ll 1$ the thermocapillary instability exists when the inequality $\alpha < 4/3$ holds and it is non–existent when this inequality is

violated. For $s_e = 1$ the static state is unstable at $\alpha < 0.692$ and stable if $\alpha > 0.692$. For the water system we calculated K_0 from (4) at $T = 373K$. Then for $d_v = d_* = 5.5 \times 10^{-6}$cm ($s_e = 1$) we get the instability criteria in terms of the accommodation coefficient: $\beta < 1.2 \times 10^{-3}$. In the light of the existing experimental data one can assume that this inequality is satisfied for real water systems where β is strongly influenced by gaseous contamination and impurity adsorption at the interface.
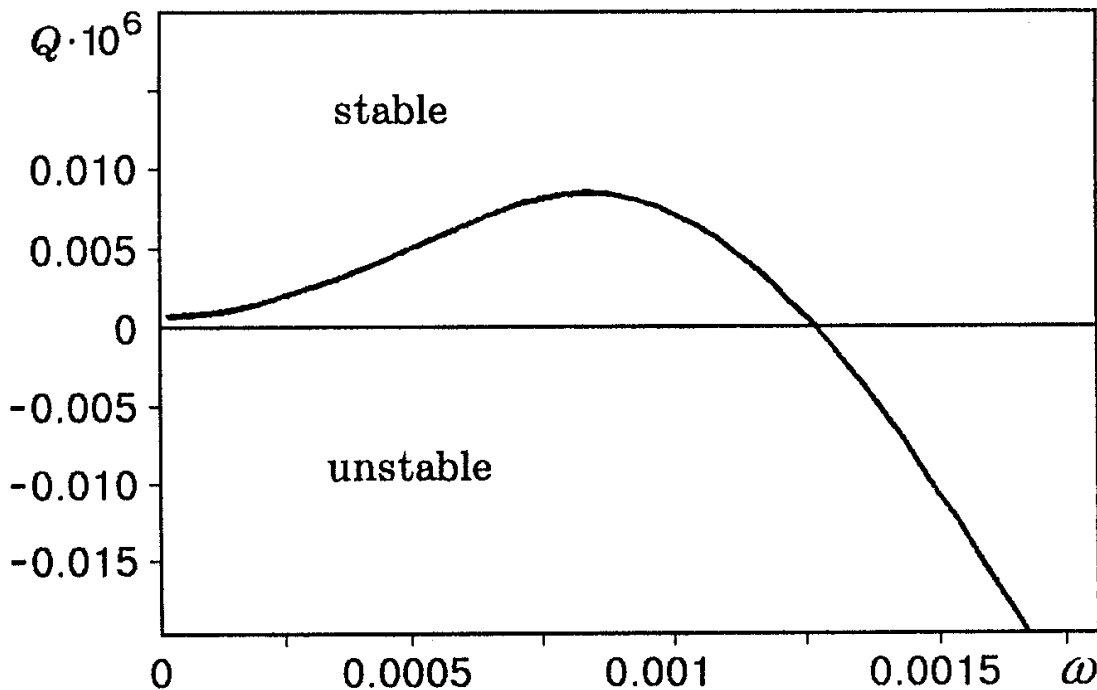

Fig. 3. Marginal stability curve for the water system at $T_0 = 373K$ in the case of a quasi-equilibrium interface, $g = 981 \times 10^{-3} cm/s^2$, $d_v = 0.05$ cm

d) In the case where it is the heat flux, rather than the temperature, that is kept constant at the heated wall, we were not able to give the explicit formulation of the thermocapillary instability criteria applicable for common liquid/vapour systems. Here the gravity influence dissapears when the parameter $G' = k_*^2 N_g/(144 N_{cr}^2)$ is much less than unity. Our initial results predict that the instability occurs at small values of the dimensionless group $\alpha' = \alpha k^*/N_{cr}$.

Note finally that all results presented in this section have been derived under the assumption of small Bond numbers. When the heated vapour layer lies above the liquid, the thermocapillary instability described above also exists and it is independent of gravity if the dimensionless parameter G (or G') is small as compared with unity.

5. On the mechanism of macrolayer formation

The model of the critical heat flux, proposed in [4] attributes this phenomenon to the evaporation of the macrolayer on a heater. In this model the macrolayer is formed on a heater after the vapour bubble detachment. The critical heat flux is attained when the heat flux is high enough to evaporate the macrolayer before the bubble detachment. The Kelvin – Helmholtz instability which is at the dividing boundary between the vapour and the superheated liquid serves to estimate the height of the macrolayer through the critical wavelength of this instability. The Kelvin – Helmholtz instability appears on the interface of two non-viscous fluids which are in relative motion with respect to each other [5]. The inviscid model for fluids excludes the possibility for the examination of the thermocapillary and phase change effects in the macrolayer. It cannot explain [9] many experimentally proven facts such as the dependence of the critical heat flux on the thickness and contamination of the heated wall.
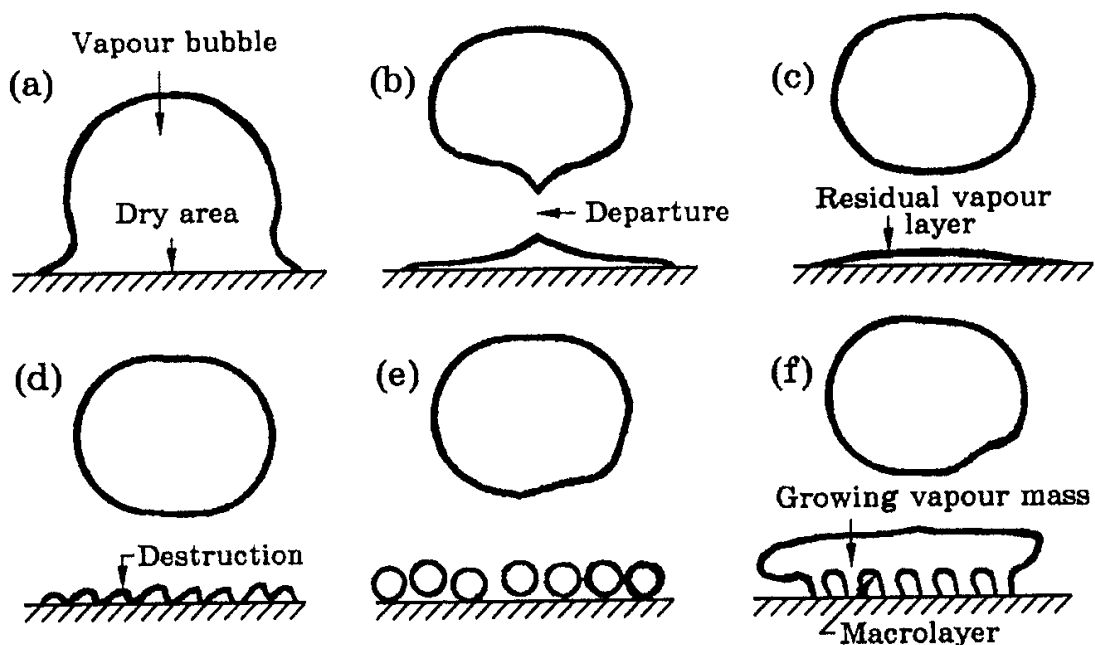

Fig. 4. Bubble departure process: (a) formation of a neck (b) break off (c) condensation of a residual vapour mass (d) destruction of a condensing vapour film (e) possible formation of small bubbles (f) mechanism of growth of the next bubble.

In a recent paper [10] the experimental observations on the macrolayer formation process are described and a new model of macrolayer formation is proposed. At high heat fluxes the observations show that when a bubble detaches and the liquid rushing onto the heater surface comes into contact with its dry area a large amount of nucleation occurs simultaneously. In the proposed model nucleate bubbles quickly grow and coalesce simultaneously producing oblate vapour film above the heated surface. The macrolayer is formed on the surface of the heater by coalescence of primary bubbles or of secondary (coalesced) bubbles.

We assume that simultaneous nucleation occurring after the liquid contact with a dry area can be explained in terms of the thermocapillary instability of a thin residual vapour layer on the heated wall condensing after the bubble break off. A possible existence of such an instability has been shown in the previous sections for sufficiently small vapour layer thickness. The departure process such as shown in Fig. 4a–c may take place (see [11]) for the vapour mass covering a dry area of a heater. When the thickness of the thinning residual vapour layer reaches a critical value the surface thermocapillary waves appear at the interface between the liquid and the residual vapour layer (Fig. 4d). These surface waves lead to the destruction of the vapour layer with the subsequent formation of a liquid macrolayer from which the vapour is supplied to the above growing vapour mass (Fig. 4f) by liquid evaporation. After complete evaporation of the macrolayer the vapour mass will grow on a dry area as shown in Fig. 4a. After the bubble detachment the macrolayer formation process can be repeated again. After the destruction of the residual vapour layer a large amount of small vapour volumes (Fig. 4d) covers the heated surface. The behavior of these small vapour bubbles (nucleate sites) cannot be predicted from the linear stability analysis performed here. They grow and then can take the spherical shape (Fig. 4e) forming the macrolayer by the coalescence. In our stability analysis the heater roughness effect is not taken into account. Thus the thermocapillary instability conception predicts that the macrolayer can be formed also on the perfectly smooth heated surfaces. This conception can be verified by experimental investigations with very smooth heater surfaces. Also, the instability analysis gives new dimensionless groups $(G, G', Q, s_n, s_e, \alpha, \alpha')$ which predict, in particular, the important role of viscous forces in the vapour phase and of the nonequilibrium phase transition inside the macrolayer. The instability conception predicts the space periodicity of the destruction process (Fig. 4d) and, as a consequence, no liquid supply into the macrolayer during its evaporation. This agrees with the existing experimental data [10]. The described instability is independent of gravity. Hence it is reasonable to assume that it also pertains to residual vapour layers on an inclined and vertical heated surfaces.

6. Conclusions

Applying an irreversible thermodynamics approach we have analyzed the possibile ways for the extension of the Hertz – Knudsen equation to moving interfaces. Knowledge of the correct rate equation is necessary to describe the thermocapillary effect in boiling where it is influenced by the mass transfer across the interface. Thermocapillary instability of a thin static vapour layer confined between a heated plate and the liquid bulk have been proven for which the critical wavelength is much less than the critical wavelength of the Rayleigh –Taylor instability. The thermocapillary instability may cause the destruction of a residual vapour layer on the heated surface leading to the formation of macrolayer beneath vapour masses.

Acknowledgements

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