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Nonlinear saturation amplitude of cylindrical Rayleigh–Taylor instability*

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The nonlinear saturation amplitude (NSA) of the fundamental mode in the classical Rayleigh–Taylor instability with a cylindrical geometry for an arbitrary Atwood number is analytically investigated by considering the nonlinear corrections up to the third order. The analytic results indicate that the effects of the initial radius of the interface (r_0) and the Atwood number (A) play an important role in the NSA of the fundamental mode. The NSA of the fundamental mode first increases gently and then decreases quickly with increasing A. For a given A, the smaller the r_0/λ (λ is the perturbation wavelength), the larger the NSA of the fundamental mode. When r_0/λ is large enough ($r_0 \gg \lambda$), the NSA of the fundamental mode is reduced to the prediction in the previous literatures within the framework of the third-order perturbation theory.

Keywords: nonlinear saturation amplitude, Rayleigh–Taylor instability, cylindrical interface

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1. Introduction

When a fluid supports another fluid of higher density in a gravity field or accelerates another fluid of higher density, the interface between the two fluids will be closely related to the Rayleigh–Taylor instability (RTI).[1,2] Assuming a heavier fluid is superposed over a lighter one in a gravitational field $-ge_y$ where g is acceleration, an initial single-mode cosine modulation, with wave number $k = 2\pi/\lambda$, where λ is the perturbation wavelength, and small perturbation amplitude ε on an interface between two fluids of densities ρ_h and ρ_l, is the simplest case. The classical linear theory[1,2] shows that the initial cosine modulation with a small amplitude grows exponentially with time t, $\eta_\mathrm{L} = \varepsilon\,\mathrm{e}^{\gamma t}$, where $\gamma = \sqrt{Akg}$ is the linear growth rate with the Atwood number $A = (\rho_\mathrm{h} - \rho_\mathrm{l})/(\rho_\mathrm{h} + \rho_\mathrm{l})$. When the typical perturbation amplitude and the wavelength are of the same order of magnitude, the second and the third harmonics are generated successively, and then the perturbation enters the nonlinear regime. Before the strong nonlinear growth regime,[3–6] there exists a weakly nonlinear growth regime.[7–24] Within the framework of the third-order weakly nonlinear theory,[7–12] the interface position at time t takes the form $\eta(x,t) = \eta_1 \cos(kx) + \eta_2 \cos(2kx) + \eta_3 \cos(3kx)$, where η_1, η_2, and η_3 are, respectively, the amplitudes of the first three harmonics

$$\eta_1 = \eta_\mathrm{L} - \frac{1}{16}\left(3A^2+1\right)k^2\eta_\mathrm{L}^3, \tag{1a}$$

$$\eta_2 = -\frac{1}{2}Ak\eta_\mathrm{L}^2, \tag{1b}$$

$$\eta_3 = \frac{1}{2}\left(A^2-\frac{1}{4}\right)k^2\eta_\mathrm{L}^3. \tag{1c}$$

For problems with a large Atwood number, $A \to 1$, equations (1a)–(1c) can be reduced to

$$\eta_1 = \eta_\mathrm{L} - \frac{1}{4}k^2\eta_\mathrm{L}^3, \tag{2a}$$

$$\eta_2 = -\frac{1}{2}k\eta_\mathrm{L}^2, \tag{2b}$$

$$\eta_3 = \frac{3}{8}k^2\eta_\mathrm{L}^3. \tag{2c}$$

As can be seen in Eq. (1a) or (2a), at the third order, the growth of the fundamental mode is reduced by the nonlinear mode-coupling effects, i.e., the third-order negative feedback to the fundamental mode. Based on the previous definition,[8,9,11,12,28] the transition into the nonlinear regime occurs when the growth of the fundamental mode is reduced by 10% in comparison to the linear growth ($\varepsilon\,\mathrm{e}^{\gamma t}$), and we have

$$\frac{\eta_\mathrm{s}}{\lambda} = \frac{\sqrt{2}}{\pi\sqrt{5(3A^2+1)}}, \tag{3}$$

where η_s is the nonlinear saturation amplitude (NSA) of the fundamental mode with corrections up to the third order. Let $A = 1$, then $\eta_\mathrm{s} \approx 0.1\lambda$ is recovered, which is a widely used threshold for nonlinearity.

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In the Cartesian geometry, the weakly nonlinear behavior of the RTI has been a field of theoretical,[7–22] experimental,[23–26] and numerical[27–32] interest. In many applications, however, the RTI occurs in a cylindrical or spherical geometry, for which the corresponding investigations are few.[33–36] The RTI plays a significant role in both astrophysics[37–39] and inertial confinement fusion (ICF).[40,41] Therefore, it is necessary to investigate the NSA of the spherical or cylindrical RTI to better understand and estimate the evolution of the RTI. In this paper, the NSA of the fundamental mode in the cylindrical RTI for irrotational, incompressible, and inviscid fluids with a discontinuous profile at an arbitrary Atwood number is investigated analytically by taking corrections up to the third order into account.

2. Theoretical framework and explicit results

This section is devoted to the detailed description of the theoretical framework of this work, and the analytic expressions of the amplitudes of the first three harmonics with corrections up to the third order are given.

A cylindrical coordinate system is established, where r and θ are normal to and along with the undisturbed interface $r = r_0$ between two fluids, respectively. The disturbed interface is located at $r = a(\theta,t)$, which is always above zero. In the following discussion, we shall denote the properties of the fluid outside the interface by subscript h and that inside the interface by subscript l, unless stated otherwise. We assume the two fluids in a gravitational field $-g e_r$ to be irrotational, incompressible, and inviscid; the governing equations for this system are

$$\frac{\partial}{\partial r}\left(r\frac{\partial \phi_i}{\partial r}\right)+\frac{\partial}{\partial \theta}\left(\frac{1}{r}\frac{\partial \phi_i}{\partial \theta}\right)=0, \text{ in two fluids,} \tag{4a}$$

$$\frac{\partial a}{\partial t}+\frac{1}{r^2}\frac{\partial a}{\partial \theta}\frac{\partial \phi_{\mathrm{l}}}{\partial \theta}-\frac{\partial \phi_{\mathrm{l}}}{\partial r}=0, \text{ at } r=a(\theta,t), \tag{4b}$$

$$\frac{\partial a}{\partial t}+\frac{1}{r^2}\frac{\partial a}{\partial \theta}\frac{\partial \phi_{\mathrm{h}}}{\partial \theta}-\frac{\partial \phi_{\mathrm{h}}}{\partial r}=0, \text{ at } r=a(\theta,t), \tag{4c}$$

$$\begin{aligned}&\rho_{\mathrm{l}}\left[\frac{\partial \phi_{\mathrm{l}}}{\partial t}+\frac{1}{2}\left(\frac{\partial \phi_{\mathrm{l}}}{\partial r}\right)^2+\frac{1}{2r^2}\left(\frac{\partial \phi_{\mathrm{l}}}{\partial \theta}\right)^2+gr\right]\\ &-\rho_{\mathrm{h}}\left[\frac{\partial \phi_{\mathrm{h}}}{\partial t}+\frac{1}{2}\left(\frac{\partial \phi_{\mathrm{h}}}{\partial r}\right)^2+\frac{1}{2r^2}\left(\frac{\partial \phi_{\mathrm{h}}}{\partial \theta}\right)^2+gr\right]\\ &+f(t)=0, \text{ at } r=a(\theta,t),\end{aligned} \tag{4d}$$

where $\phi_i(r,\theta,t)$ are the velocity potentials for the two fluids with i being h or l, $f(t)$ is an arbitrary function of time t, which is introduced in the integral for the Bernoulli relation with respect to space, and the perturbation interface $a(\theta,t)$ corresponds to $\eta(x,t)$ in the Cartesian geometry. The Laplace equation (4a) comes from the incompressibility condition in the cylindrical geometry. Equations (4b) and (4c) represent the kinematic boundary conditions in the cylindrical geometry (the normal velocity continuous condition at the interface), i.e., a fluid particle initially situated at the material interface remains at the interface afterwards. The Bernoulli equation (4d) represents the dynamic boundary condition, i.e., the pressure continues across the material interface.

We consider an initial perturbation in the form $r = a(\theta,t=0) = r_0 + \varepsilon\cos(\kappa\theta)$, where r_0 is a positive constant, mode number $\kappa = 2\pi r_0/\lambda$, and $\varepsilon \ll \lambda$. Due to this small-amplitude perturbation in the cylindrical interface, the perturbed interface is prone to the RTI. Higher harmonics (i.e., the second harmonic, the third harmonic, and so on) will be subsequently generated in the nonlinear mode-coupling process. Hence, $a(\theta,t)$ and $\phi_i(r,\theta,t)$ can be expanded into a power series in $\hat{\varepsilon} = \varepsilon/\lambda$ (the initial perturbation amplitude normalized by the perturbation wavelength, which is a small parameter) as

$$\begin{aligned}a(\theta,t) &= \zeta(t)r_0+\sum_{n=1}^{N}a^{(n)}(\theta,t)\\ &= r_0\left(1+\sum_{n=1}^{[N/2]}\hat{\varepsilon}^{2n}\lambda^{2n}\mathrm{e}^{2n\beta t}\alpha_{2n,0}\right)\\ &\quad+\sum_{n=1}^{N}\hat{\varepsilon}^{n}\lambda^{n}\mathrm{e}^{n\beta t}\sum_{m=0}^{[n/2]-1}\alpha_{n,n-2m}\\ &\quad\times\cos(n-2m)\kappa\theta+\mathrm{O}(\hat{\varepsilon}^{N+1}),\end{aligned} \tag{5a}$$

$$\begin{aligned}\phi_{\mathrm{l}}(r,\theta,t) &= \sum_{n=1}^{N}\phi_{\mathrm{l}}^{(n)}(r,\theta,t)\\ &= \sum_{n=1}^{N}\hat{\varepsilon}^{n}\lambda^{n}\mathrm{e}^{n\beta t}\sum_{m=0}^{[n/2]}\phi_{l,n,n-2m}r^{(n-2m)\kappa}\\ &\quad\times\cos(n-2m)\kappa\theta+\mathrm{O}(\hat{\varepsilon}^{N+1}),\end{aligned} \tag{5b}$$

$$\begin{aligned}\phi_{\mathrm{h}}(r,\theta,t) &= \sum_{n=1}^{N}\phi_{\mathrm{h}}^{(n)}(r,\theta,t)\\ &= \sum_{n=1}^{N}\hat{\varepsilon}^{n}\lambda^{n}\mathrm{e}^{n\beta t}\sum_{m=0}^{[n/2]}\phi_{h,n,n-2m}r^{-(n-2m)\kappa}\\ &\quad\times\cos(n-2m)\kappa\theta+\mathrm{O}(\hat{\varepsilon}^{N+1}),\end{aligned} \tag{5c}$$

where function $\zeta(t)$ determines whether the unperturbed interface moves with time, the interface will keep resting when $\zeta(t) \equiv 1$, otherwise, it will move away from the initial position $r(t=0) = r_0$. The $a^{(n)}(\theta,t)$ and $\phi_{\mathrm{l}}^{(n)}(r,\theta,t)$ ($\phi_{\mathrm{h}}^{(n)}(r,\theta,t)$) are, respectively, the n-th order perturbed interface and the n-th order perturbed velocity potential for the inner (outer) fluid of the interface when the first three harmonics are taken into account. For the $(n-2m)$-th Fourier harmonic at the n-th order, when $m = 0$, $a_{n-2m}^{(n)} = \hat{\varepsilon}^{n}\lambda^{n}\mathrm{e}^{n\beta t}\alpha_{n,n-2m}$ is the generation coefficient of the perturbation interface, $\phi_{\mathrm{l},n-2m}^{(n)} = \hat{\varepsilon}^{n}\lambda^{n}\mathrm{e}^{n\beta t}\phi_{\mathrm{l},n,n-2m}r^{(n-2m)\kappa}$ and $\phi_{\mathrm{h},n-2m}^{(n)} = \hat{\varepsilon}^{n}\lambda^{n}\mathrm{e}^{n\beta t}\phi_{\mathrm{h},n,n-2m}r^{-(n-2m)\kappa}$ are the generation coefficients of the velocity potentials for the light and the heavy fluids, respectively; when $m > 0$, they are the corresponding correction coefficients of the n-th order for the perturbation interface. Here Gauss's symbol $[n/2]$ denotes the maximum integer that is less than or equal to $n/2$, and β is the linear growth rate in the cylindrical geometry. Note that the perturbation velocity potentials $\phi_{\rm h}(r,\theta,t)$ and $\phi_{\rm l}(r,\theta,t)$ have satisfied the Laplace equation (4a) and the boundary conditions $\nabla\phi_{\rm h}|_{r\to+\infty}=0$ and $\nabla\phi_{\rm l}|_{r=0}=0$. Also, $\alpha_{1,1}=1$ when the initial condition is taken into account. The coupling factors in the amplitudes of the Fourier harmonics, $\alpha_{n,n-2m}$ ($n=2,\ \ldots,\ N$, $m=0,\ 1,\ldots,\ [n/2]$), and β are to be determined.

It should be emphasized that the procedure to solve this system is non-trivial. The detailed steps are as follows. (i) Substitute Eqs. (5a)–(5c), in which $\mathrm{O}(\hat{\varepsilon}^{N+1})$ is neglected, into Eqs. (4b)–(4d). (ii) Replace r in the three equations with $a(t)$. (iii) Re-express the left-hand sides of the resulting equations into Maclaurin series of $\hat{\varepsilon}$ and collect terms of the same power in $\hat{\varepsilon}$ to construct a set of equations. (iv) Eliminate unknown factors in the velocity potentials and solve the resulting equations successively for $n=1,2,\ldots,N$.

By employing the above steps, the linear growth rate and the coupling factors of the first three harmonics with corrections up to the third order (i.e., $N=3$ in Eqs. (5a)–(5c)) can be expressed as

$$\beta=\sqrt{\frac{Ag\kappa}{r_0}}, \tag{6a}$$

$$\alpha_{2,0}=-\frac{1}{4r_0^2}, \tag{6b}$$

$$\alpha_{2,2}=-\frac{A\kappa+1}{2r_0}, \tag{6c}$$

$$\alpha_{3,1}=\frac{-3A^2\kappa^2+A\kappa-\kappa^2+9}{16r_0^2}, \tag{6d}$$

$$\alpha_{3,3}=\frac{4A^2\kappa^2+7A\kappa-\kappa^2+3}{8r_0^2}. \tag{6e}$$

Expression (6a) shows that the linear growth rates in cylindrical and Cartesian geometries are different unless $\kappa/r_0=k$ (i.e., the same λ). Keeping Atwood number A, acceleration g, and mode number κ fixed, the smaller the initial radius of the interface (r_0), the larger the linear growth rate in the cylindrical geometry. Expression (6c) denotes that the second harmonic has a character of negative growth (i.e., anti-phase). In addition, expressions (6c)–(6e) demonstrate that the coupling factors are influenced not only by A but also by κ and r_0. If the constant λ is considered in both the cylindrical and the Cartesian geometries (i.e., $\kappa/r_0=k$), and r_0 is large (i.e., $r_0\to+\infty$), $\alpha_{n,n-2m}/k^{n-1}$ ($n=2,\ldots,N,m=0,1,\ldots,\ [n/2]$) will be simplified to the corresponding Atwood number $f_{n,n-2m}$ in Ref. [11]. This means that under the conditions of the same λ and a large r_0, the perturbed interface in the cylindrical geometry reproduces that in the Cartesian geometry, and the results from the classical third-order weakly nonlinear theory,[7–11] as shown in Eqs. (1a)–(1c), are recovered. It should be noted that the generation of $\alpha_{2,0}$ is an essential character different from the results in the Cartesian geometry where $\alpha_{2,0}=0$.

Accordingly, the interface position in the framework of the third-order theory in the cylindrical geometry takes the form $\zeta r_0+a(\theta,t)\doteq\zeta r_0+\sum_{n=1}^{3}a_n\cos(n\kappa\theta)$, where ζ and the amplitude of the n-th harmonic, a_n, are

$$\zeta = 1+\eta_{\rm Lc}^2\alpha_{2,0}, \tag{7a}$$

$$a_1 = \eta_{\rm Lc}\left(1+\eta_{\rm Lc}^2\alpha_{3,1}\right), \tag{7b}$$

$$a_2 = \eta_{\rm Lc}^2\alpha_{2,2}, \tag{7c}$$

$$a_3 = \eta_{\rm Lc}^3\alpha_{3,3}, \tag{7d}$$

where $\eta_{\rm Lc}=\varepsilon\,\mathrm{e}^{\beta t}$ is the linear growth amplitude of the fundamental mode in the cylindrical geometry. It should be pointed out that the amplitude of the fundamental mode is just corrected by the third harmonic. An essential character different from the Cartesian RTI is that the zeroth order harmonic does not vanish in the cylindrical RTI (see Eq. (7a)). This means that the position of the initial unperturbed interface $r=r_0$ will be changed into $r=\zeta(t)r_0$ with the development of the perturbation, entirely different from that in the Cartesian space where the initial unperturbed interface stays invariable all the time.

3. NSA of the fundamental mode

We start to analyze the NSA of the fundamental mode with corrections up to the third order. As mentioned above, the NSA of the fundamental mode can be defined as the linear growth amplitude of the fundamental mode $\left(\sim \mathrm{e}^{\beta t}\right)$ at the saturation time ($t_{\rm s}$) when the growth of the fundamental mode is reduced by 10% in comparison to the linear growth. Hence, we have

$$\frac{\varepsilon\,\mathrm{e}^{\beta t_{\rm s}}-a_1(t_{\rm s})}{\varepsilon\,\mathrm{e}^{\beta t_{\rm s}}}=\frac{1}{10}. \tag{8}$$

Here, β and a_1 are, respectively, substituted by Eqs. (6a) and (7b), and the saturation time $t_{\rm s}$ of the fundamental mode can be obtained by solving the resulting Eq. (8) as

$$t_{\rm s}=\sqrt{\frac{r_0}{Ag\kappa}}\log\left(\frac{2r_0}{\varepsilon}\frac{\sqrt{2/5}}{\sqrt{3A^2\kappa^2-A\kappa+\kappa^2-9}}\right). \tag{9}$$

Then, the NSA of the fundamental mode with corrections up to the third order is determined as $a_{\rm s}=\eta_{\rm Lc}(t_{\rm s})=\varepsilon\,\mathrm{e}^{\beta t_{\rm s}}$, i.e.,

$$a_{\rm s}=2\sqrt{\frac{2}{5}}\sqrt{\frac{r_0^2}{3A^2\kappa^2-A\kappa+\kappa^2-9}}. \tag{10}$$

By using mode number $\kappa=2\pi r_0/\lambda$, the normalized saturation time and the NSA of the fundamental mode are

$$\frac{t_{\rm s}}{\sqrt{\lambda/g}}=\frac{\log 2\sqrt{\dfrac{2r_0^2}{5\varepsilon\left(12\pi^2A^2r_0^2-2\pi Ar_0+4\pi^2r_0^2-9\right)}}}{\sqrt{2A\pi}}, \tag{11a}$$

$$\frac{a_{\rm s}}{\lambda}=\frac{2\sqrt{2}r_0}{\sqrt{5}\sqrt{12\pi^2A^2r_0^2-2\pi Ar_0+4\pi^2r_0^2-9}}. \tag{11b}$$

When the limit of $r_0 \to +\infty$ is taken into account, we have

$$\frac{t_{\mathrm{s}}^{\infty}}{\sqrt{\lambda/g}} = \frac{\log\sqrt{2/5\varepsilon\pi^2(3A^2+1)}}{\sqrt{2A\pi}}, \tag{12a}$$

$$\frac{a_{\mathrm{s}}^{\infty}}{\lambda} = \frac{\sqrt{2}}{\pi\sqrt{5}\sqrt{3A^2+1}}. \tag{12b}$$

As can be seen, when r_0/λ is large enough (i.e., $r_0/\lambda \to +\infty$), the NSA of the fundamental mode in the cylindrical geometry will tend to that in the Cartesian geometry, i.e., equation (3) is recovered. This means that the cylindrical effect on the NSA of the fundamental mode will vanish.

We show the normalized NSA, a_{s}/λ, of the fundamental mode versus Atwood number A for different initial radii of the interface in Fig. 1. Here r_0/λ =0.8, 1.0, 2.0, $+\infty$ are uniformly selected.

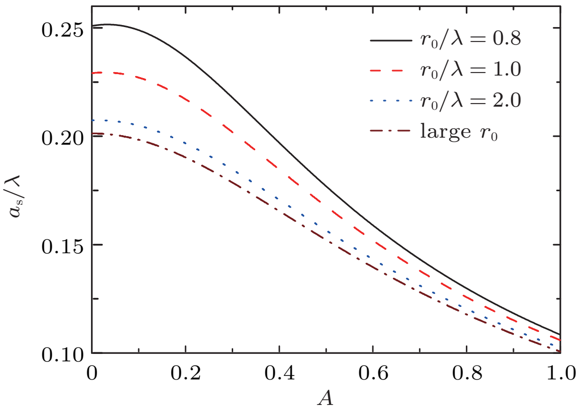

Fig. 1. (color online) The normalized NSA of the fundamental mode, a_{s}/λ, versus arbitrary A for different initial radii of the cylindrical interface. The result of large r_0 corresponds to that of the Cartesian geometry.

Figure 1 demonstrates that the normalized NSA of the fundamental mode decreases monotonously with increasing A for large r_0 (corresponding to the Cartesian geometry), while it firstly rises weakly to a peak and then decreases sharply with increasing A for finite r_0/λ. Thus, there is a critical A_{c}. For $A \leqslant A_{\mathrm{c}}$, the NSA keeps increasing to a maximum with increasing A; otherwise, it keeps decreasing. Meanwhile, A_{c} has a trend of increasing with decreasing r_0/λ. For a selected A, the NSA of the fundamental mode increases with decreasing r_0/λ. It is clear that r_0 has a significant influence on the NSA of the fundamental mode, particularly when $r_0 < \lambda$. As can be seen, the normalized NSA of the fundamental mode at $A = 1$ and $r_0/\lambda = 0.8$ is 0.111, while that at $A = 1$ for large r_0 is only 0.1. Furthermore, at $A = 0$ and $r_0/\lambda = 0.8$, the NSA of the fundamental mode is 0.275, while the corresponding result in the Cartesian geometry is just 0.201. Accordingly, the cylindrical effects (especially at small r_0/λ) play an important role in the NSA of the fundamental mode for arbitrary A.

4. Conclusion

The nonlinear saturation amplitude of the fundamental mode in the classical RTI (irrotational, incompressible, and inviscid fluids) with a discontinuous cylindrical profile for an arbitrary Atwood number and nonlinear corrections up to the third order is explored analytically. The prediction of the NSA of the fundamental mode with corrections up to the third order from the classical weakly nonlinear theory[7–11] is recovered when the initial radius of the interface normalized by the perturbation wavelength tends to infinity. The NSA of the fundamental mode in the cylindrical geometry has a different trend with Atwood number A from the one in the Cartesian geometry. The NSA of the fundamental mode in the Cartesian geometry decreases monotonously with increasing A; when the cylindrical effects are taken into account, it first slightly increases to a maximum and then decreases with increasing A. Accordingly, there is a critical A_{c} which is small. When $A \leqslant A_{\mathrm{c}}$, the NSA keeps increasing generally with increasing A; when $A > A_{\mathrm{c}}$, the NSA keeps decreasing quickly with increasing A. It is also found that the smaller the normalized initial radius of the interface r_0/λ, the larger A_{c}. Again, r_0 plays a vital role in the NSA of the fundamental mode. The NSA of the fundamental mode increases with decreasing r_0/λ, especially when r_0 is compared to λ. Our analytic results show that not only the Atwood number but also the initial radius of the interface strikingly influences the NSA of the RTI. Thus, it should be included in applications where the NSA plays a role, such as the inertial confinement fusion ignition target design.

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