Nonlinear Electrohydrodynamic Stability of Two Superposed Bounded Fluids in the Presence of Interfacial Surface Charges

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The method of multiple scales is used to analyse the nonlinear propagation of waves on the interface between two superposed dielectric fluids with uniform depths in the presence of a normal electric field, taking into account the interfacial surface charges. The evolution of the amplitude for travelling waves is governed by a nonlinear Schrödinger equation which gives the criterion for modulational instability. Numerical results are given in graphical form, and some limiting cases are recovered. Three cases, in the pure hydrodynamical case, depending on whether the depth of the lower fluid is equal to or greater than or smaller than the one of the upper fluid are considered, and the effect of the electric field on the stability regions is determined. It is found that the effect of the electric field is the same in all the cases for small values of the field, and there is a value of the electric field after which the effect differs from case to case. It is also found that the effect of the electric field is stronger in the case where the depth of the lower fluid is larger than the one of the upper fluid. On the other hand, the evolution of the amplitude for standing waves near the cut-off wavenumber is governed by another type of nonlinear Schrödinger equation with the roles of time and space are interchanged. This equation makes it possible to determine the nonlinear dispersion relation, and the nonlinear effect on the cut-off wavenumber.

Key words: Hydrodynamic Stability; Electrohydrodynamics; Nonlinearity; Interfacial Instability; Dielectric Fluids; Surface Charges.

1. Introduction

The two-dimensional evolution of a nonlinear wave packet propagating on deep water has been investigated by Lighthill [1], Whitham [2], and Yuen and Lake [3], using the method of averaged Lagrangian. Also Chu and Mei [4], Hasimoto and Ono [5] used the multiple scales method to study the same problem. All the above authors derived two-dimensional nonlinear Schrödinger equation describing the modulation of the wave amplitude. Zakharov [6] showed that this equation provides an elegant approach to examine the modulational instability of finite amplitude waves. It was shown by Yuen and Lake [3] that the nonlinear Schrödinger equation can be derived by the averaged Lagrangian method when the spatial variations in the amplitude are included in the dispersion relation. Moreover, they demonstrated that the two-dimensional nonlinear Schrödinger equation provides a quantitative satisfactory description of the long-time evolution of weakly nonlinear wave packets.

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The nonlinear modulation of waves propagating along the interface between two liquid layers has been investigated by Qi-su [7] and Tanaka [8]. By using the method of multiple scales, the evolution equation of a wave packet of the wave train has been found. Then they discussed the stability of a wave train with infinitesimal perturbation in the direction of propagation of the wave train and obtained the stability criterion. But their theory can not be used to discuss the instability when transverse perturbation occurs. For recent works concerning the instability of capillary-gravity waves, see the work of Chhabra and Khosla [9], Christodoulides and Dias [10], Collin et al. [11], Jones [12], Dolai [13], Qingpu [14], Lee [15], and Kato et al. [16].

On the other hand, electrohydrodynamics is the field of the mechanics of continua that studies the motion of media interacting with the electric field. Such an interaction takes place as a result of action of the Coulomb force upon a medium, or as a result of work of the electric field in flowing of currents. The motion of a medium gives rise to re-distribution of a volume charge, which results in changing the electric field and, hence, the force acting on a medium. In the majority of problems under consideration, electric fields or electric charges are specified by external sources. Such a situation takes place
during operation of electrohydrodynamic generators, pumps, separators, filters, and other devices. This subject is treated in a vast literature. The last studied is the class of electrohydrodynamic problems in which the electric field or electric charges arise as a result of a contact between media of different nature: liquid-solid body, liquid-gas or two different liquids (see [17]–[22]).

The effect of the electric field on the motion of fluids has been studied by a number of scientists since the pioneering work of Rayleigh [23], Stokes [24], and Melcher [25]. Michael [26] investigated the stability of an incompressible, inviscid, perfectly conducting fluid layer in the presence of electrostatic forces, and he found that these forces can have a destabilising effect on the fluid motions. Shivamoggi [27] has also examined the stability of such a problem in the neighbourhood of the linear cut-off wavenumber. Kant et al. [28] investigated the stability of weakly nonlinear waves on the surface of a perfectly conducting fluid layer in the presence of an applied electric field by using the derivative expansion method. The nonlinear electrohydrodynamic Rayleigh-Taylor instability was investigated by Mohamed and Elshehawey [29]. They obtained two nonlinear Schrödinger equations by means of which one can deduce the cut-off wavenumber. Quite recently, Elshehawey [30] investigated the same problem of Rayleigh-Taylor instability for a normal electric field normal to the interface $\mathbf{E}_0$ where the superscripts 1 and 2 refer to the lower and upper fluids, respectively.

If the two fluids are assumed to be inviscid and incompressible, and the fluid motion being irrotational, then there exist velocity potentials $\phi^{(j)}(x, y, t)$ within the two regions such that $\mathbf{V}^{(j)} = \nabla \phi^{(j)}$. Since the system is stressed by a normal electric field, we shall assume that it allows for the presence of surface charges at the interface such that $\mathbf{E}^{(1)} \mathbf{E}_0^{(2)} \neq \mathbf{E}^{(2)} \mathbf{E}_0^{(1)}$.

We shall assume that the quasi-static approximation is valid and the electric field $\mathbf{E}$ is irrotational. The electric potentials $\psi^{(j)}$ are defined such that

$$E^{(j)} = -E_0^{(j)} e_y - \nabla \psi^{(j)},$$

$$E_0^{(j)} = \frac{V_0}{\varepsilon^{(j)}} = \frac{V_0}{\varepsilon^{(j)}} e_y = \frac{V_0}{\varepsilon^{(j)}} e_y,$$

where $e_y$ is the unit vector in the y-direction.

The basic equations relevant to our problem, are

$$\phi_{xx}^{(j)} + \phi_{yy}^{(j)} = 0 \quad \text{and} \quad \psi_{xx}^{(j)} + \psi_{yy}^{(j)} = 0,$$

$$j = 1, 2,$$

where $j = 1, 2$ represent the regions $-h < y < \eta(x, t)$ and $\eta(x, t) < y < b_0$, respectively, $y = \eta(x, t)$ is the elevation of the interface measured from the unperturbed level, and $t$ denotes the time.

The various physical quantities are normalized with respect to a characteristic length $l_c = (T/\rho_1 g)^{1/2}$ and a characteristic time $t_c = (l_c^2 g)^{1/2}$ and the characteristic potential functions $\phi = (g l_c^2)^{1/2}$, $\psi = (\rho_1 g l_c^2)^{1/2}$, where $g$ is the acceleration due to gravity acting in the negative y-direction and $T$ is the surface tension. Hence we have $l = l_c$, $t = t_c$, $\phi^{(j)} = \phi^{(j)}$, $\psi^{(j)} = \psi^{(j)}$, and $\eta(x, t)$ is the elevation of the interface measured from the unperturbed level, and $t$ denotes the time.

The boundary conditions at the interface $y = \eta(x, t)$ are

$$\eta_t - \phi_x^{(j)} + \eta_x \phi_x^{(j)} = 0, \quad j = 1, 2,$$

$$\eta_x \langle \psi_x \rangle + \langle \psi_x \rangle = \eta_x \langle E_0 \rangle,$$

$$\eta_x \langle \psi_x \rangle - \langle \psi_y \rangle = 0,$$

where $\langle \cdot \rangle$ denotes the average over the fluid layer.

2. Formulation of the Problem

We consider the two-dimensional wave motion on the interface $y = 0$ between the two superposed dielectric fluids with uniform depths, the upper fluid having the density $\rho^{(2)}$, dielectric constant $\varepsilon^{(2)}$ and being bounded by the conducting plane $y = b_0$ which is raised to the potential $V_0$ whereas the lower fluid with density $\rho^{(1)}$, dielectric constant $\varepsilon^{(1)}$ is bounded from below by an earthed conducting plane $y = -a_0$. As a result of the potential difference between the planes, both fluids are subjected to a constant electric field normal to the interface $E_0^{(2)}$ and $E_0^{(1)}$, where the superscripts 1 and 2 refer to the lower and upper fluids, respectively.
\[
\phi_i^{(1)} - \rho \phi_i^{(2)} + (1 - \rho) \eta + \frac{1}{2} \left( \phi_x^{(1)} - \rho \phi_x^{(2)} \right)^2 + \frac{1}{2} \left( \phi_y^{(1)} - \rho \phi_y^{(2)} \right)^2 = \\
\eta_{xx} \left( 1 + \eta_{y}^2 \right) \frac{3}{2} - \frac{1}{2} \left( \langle \psi_x \psi_x \rangle \right) + \frac{1}{2} \left( \langle \psi_y \psi_y \rangle \right) - \left( \langle \varepsilon \psi_E \psi_y \rangle \right) + 2 \eta_x \left( \langle \varepsilon \psi_E \psi_x \rangle \right) - 2 \eta_y \left( \langle \varepsilon \psi_x \psi_y \rangle \right) \\
- \eta_{x}^2 \left( \langle \varepsilon \psi_E^2 \rangle \right) + 2 \eta_{y}^2 \left( \langle \varepsilon \psi_E \psi_y \rangle \right) + \text{higher order terms},
\]

where \( \rho = \rho^{(2)} / \rho^{(1)} \), and \( \langle \cdot \rangle \) represents the jump across the interface.

The solutions \( \phi^{(j)}, \psi^{(j)} \) of (2) must satisfy the following conditions at the boundaries \( y = -a, b \):

\[
\phi_y^{(j)}, \psi_y^{(j)} = 0 \quad \text{on} \quad y = -a, b,
\]

where \( a \) and \( b \) are the dimensionless quantities for the depths.

To investigate the modulation of a weakly nonlinear wave with narrow band width spectrum, we employ the method of multiple scales \([37]\) by introducing the variables 
\[
x_n = e^{\gamma x} \quad \text{and} \quad t_n = e^{\gamma t}, \quad (n = 0, 1, 2, 3),
\]

and expanding \( \eta, \phi^{(j)} \) and \( \psi^{(j)} \), \( j = 1, 2 \) in the asymptotic series

\[
F(x,y,t) = \sum_{n=1}^{3} \hat{e}^n F_n \left( x_0, x_1, \ldots, x_N, y, t_0, t_1, \ldots, t_N \right) + O(\hat{e}^4),
\]

where the small parameter \( \hat{e} \) characterizes the steepness ratio of the wave, and the expansion of \( \eta \) according to (9) is independent of \( y \).

Expanding now the boundary conditions (3)-(7) into Maclaurin series expansions around \( y = 0 \), then substituting (8) into (2) and the boundary conditions (3)-(7), and equating the coefficients of equal powers in \( \hat{e} \), we obtain the linear and successive nonlinear partial differential equations for \( \eta, \phi, \psi^{(j)} \) \([36]\); they will not be given because they are very lengthy.

### 3. Linear Theory

We assume that there is no steady flow in the undisturbed state, so that we choose the following quasi-monochromatic wave as the starting solutions to the first order problem

\[
\eta_1 = i \sigma_1 A e^{i\theta} + \text{c.c.},
\]

\[
\phi_1^{(1)} = \frac{\omega \cosh(\kappa y + a')}{k \cosh a'} A e^{i\theta} + \text{c.c.} + B_1^{(1)},
\]

where \( a' = k a, b' = k b, \sigma_1 = \tanh a', \sigma_2 = \tanh b' \), and \( \theta (= k \omega_0 - \alpha \omega_0) \) is the phase of the carrier wave, \( k \) and \( \omega \) being, respectively, the wavenumber and frequency of the centre of the wave packets, and c.c. stands for the complex conjugate of the preceding term (or terms), and \( i \) is the imaginary unit. Here, the complex amplitude \( A \) and the additional real constants \( B_1^{(1)} \) (which represent the arbitrariness associated with the velocity potential) are functions of the slow scales \( x_1, x_2, t_1, t_2 \).

In order that the starting solution should not be trivial, the wavenumber \( k \) and the frequency \( \omega \) must satisfy the dispersion relation

\[
\omega^2 = k \left( \frac{1}{\sigma_1} + \frac{\rho}{\sigma_2} \right)^{-1} \cdot \left[ 1 - \rho + k^2 - k \left( \frac{V_2^2}{\sigma_2} + \frac{V_1^2}{\sigma_1} \right) \right],
\]

where

\[
V_j = e^{(j)(-1)} E_0^{(j)} \quad \text{for} \quad j = 1, 2.
\]

The dispersion relation (14) was initially obtained by Melcher \([25]\), Mohamed and Elshehawey \([29]\) (for the case of two-dimensional, semi-infinite fluids), and also by Qi-Su \([7]\) (for the corresponding hydrodynamical case, i.e. when \( E_0^{(j)} = 0 \) or \( e^{(1)} = e^{(2)} \)). The critical wavenumber \( k_c \) at which \( \omega = 0 \) is the linear electrohydrodynamic cut-off wavenumber separating stable from unstable disturbances. For a numerical discussion of the dispersion equation (14), we have a transition curve, namely

\[
k^2 - k \left( \frac{V_2^2}{\sigma_2} + \frac{V_1^2}{\sigma_1} \right) + (1 - \rho) = 0.
\]
From (15) we calculated the values of $k$ corresponding to some different values of $V_1^2$ and $V_2^2$ in the cases when $a' = b'$, $a' > b'$ and $a' < b'$, respectively, for values of the density parameter $p \geq 1$. We plotted the neutral stability condition (15), separating the stable $S$ and unstable $U$ regions. Figures 1 and 2 are the stability diagrams for the linear case in the $k - \rho$ plane for different values of the electric field variations.

Figure 1 a is drawn for the case of no electric field influence, i.e. $V_1^2 = V_2^2 = 0$ for any depths of the fluids. The resulting curve represent the neutral curve separating the stable and unstable regions. Figure 1 b is drawn for the case of equal depths (i.e. $a' = b' = 0.9$), and the resulting curves correspond to the electric field values $V_2^2 = 0.05, 0.1, 0.15, 0.2$ and 0.25, respectively, with $V_1^2 = 0.009$. It is clear from Fig. 1 a that increasing the electric field values increases the unstable region, i.e. the electric field has a destabilizing effect.

Figure 2 a is drawn for the case when the lower depth is greater than the upper one (i.e. $a' = 0.9$ and $b' = 0.3$), while Fig. 2 b is drawn for the case when the lower depth is less than the upper one (i.e. $a' < b'$, where $a' = 0.3$ and $b' = 0.9$). The resulting curves in Fig. 2 correspond to the same values of the electric field mentioned in Figure 1 b. It is clear from Fig. 2 that the normal electric field has usually a destabilizing effect and that the effect is stronger or faster in the case $a' > b'$ than the other cases when $a' < b'$ and $a' = b'$.

Now, to derive the equation for the evolution of travelling waves, we need to proceed to the second order and higher order problems.

4. Second Order Solution

Since our aim is to study the amplitude modulation for travelling waves when $\phi^2 > 0$, we now proceed to the second order problem in $O(\varepsilon^2)$. With the use of the first order solutions given by (9)–(13) on the right-hand sides of the second order equations, and solving the resulting equations, the second order solutions $\eta_2$, $\phi_2^{(1)}$ and $\psi_2^{(1)}$ ($j = 1, 2$), take the form

\[
\eta_2 = \left(\frac{a' + \sigma_1}{k} \frac{\partial A}{\partial x_1} + \frac{\sigma_1}{\omega} \frac{\partial A}{\partial t_1} \right) e^{i\theta} + AA^2 e^{2i\theta} + c.c. + \xi_2,
\]

\[
\phi_2^{(1)} = -\frac{i\omega (ky + a') \sinh (ky + a')}{k^2 \cosh a'} \frac{\partial A}{\partial x_1} e^{i\theta} - \frac{i\omega (1 + \sigma_1^2)}{2k\sigma_1} \left(\Lambda + k\sigma_1\right) \frac{\cosh (ky + a')}{\cosh 2a'} A^2 e^{2i\theta} + c.c. + B_2^{(1)},
\]

\[
\phi_2^{(2)} = \frac{i\omega \sigma_1}{k^2 \sigma_2} \left[\frac{(ky - b') \sinh (ky - b')}{\cosh b'} + \left(\frac{a' - a'}{\sigma_1 - \sigma_2}\right) \frac{\cosh (ky - b')}{\cosh b'}\right] \frac{\partial A}{\partial x_1} e^{i\theta}
+ \frac{i\omega (1 + \sigma_1^2)}{2k\sigma_2} \left(\Lambda - k\sigma_1\right) \frac{\cosh 2(ky - b')}{\cosh 2b'} A^2 e^{2i\theta} + c.c. + B_2^{(2)},
\]

\[
\psi_2^{(1)} = E_0^{(1)} \left[\left\{\frac{\sinh (ky + a') + (ky + a') \cosh (ky + a')}{k \cosh a'}\right\} \frac{\partial A}{\partial x_1} + \frac{\sinh (ky + a')}{\omega \cosh a'} \frac{\partial A}{\partial t_1} \right] e^{i\theta}
+ \frac{E_0^{(1)} (1 + \sigma_1^2)}{2\sigma_1} \left(\Lambda + k\sigma_1\right) \frac{\sinh 2(ky + a')}{\cosh 2a'} A^2 e^{2i\theta} + c.c.,
\]
where $\xi_2$ (which represents the induced mean motion or the zero frequency correction to slow modulation of the fundamental mode), and $B^{(j)}_V$ are real functions of the slow scales $x_1, x_2, t_1,$ and $t_2$ to be determined by considering the equations of higher orders, and

$$\Lambda = \frac{1}{G_1} \left[ \frac{1}{2} \omega^2 \sigma_1^2 (1 - \rho) - 3 \frac{\omega^2}{2} - 2 \left( 1 - \frac{\rho \sigma_1^2}{\sigma_2^2} \right) \right]$$

(21)

On substitution from (16)–(20) into the last condition of the second order equations, we obtain the non-secularity condition, which consists of two parts; one is

$$\frac{\partial A}{\partial t_1} + v_g \frac{\partial A}{\partial x_1} = 0$$

(23)

together with its c.c., where $v_g = \frac{d\omega}{dk}$ is the group velocity of the wave train, and the other one is

$$\frac{(1 - \rho) \xi_2 + \frac{\partial B^{(1)}_1}{\partial t_1} - \rho \frac{\partial B^{(2)}_1}{\partial t_1} + G_2 |A|^2 = 0}$$

(24)

where

$$G_1 = \frac{\omega^2}{k} \left( \sigma_1 + \rho \sigma_2 \right) - 3 k^2 + k \left( V_2^2 \sigma_2 + V_1^2 \sigma_1 \right)$$

(22)

and

$$S_j^2 = \sigma_j^2 (1 - \sigma_j^2), \quad j = 1, 2.$$

Equation (23) implies that, to the lowest order in $\xi$, the complex amplitude $A$ remains constant in a frame of reference moving with the group velocity $v_g$ of the waves, that is, $A$ depends on $x_1$ and $t_1$ only through $\xi$ defined as $\xi = x_1 - v_g t_1 = \xi(x - v_g t)$. In (21), the case when $G_1 = 0$, for which $\eta_2, \phi^{(j)}_2$ and $\psi^{(j)}_2$ become infinite, corresponds to the case of second harmonic resonance which can be dealt with along the lines outlined by Singla et al. [20], in another problem of interest. Such a kind of resonance in our case will be discussed separately in another paper in the near future. In this section, we have assumed that this quantity is different from zero in (16)–(20).

5. Third Order Solution

Let us proceed to the third order problem. Introducing (9)–(13) and (16)–(20) into the third order perturbation equations and solving the resulting equations, we obtain the third order solutions $\eta_3, \phi^{(j)}_3$ and $\psi^{(j)}_3$ as indicated in the Appendix. On substitution from (A.1)–(A.5) into the last condition of the third order problem, we obtain the non-secularity condition from the coefficient of $e^{i\theta}$, that is

$$\psi^{(2)}_2 = - \frac{E^{(2)}_0}{\sigma_2} \left[ \left( \frac{a' - b'}{\sigma_1 \sigma_2} \right) \frac{\sinh(ky - b')}{k \cosh b'} + \frac{(ky - b') \cosh(ky - b')}{k \cosh b'} \right] \frac{\partial A}{\partial x_1} + \frac{\sinh(ky - b')}{\omega \cosh b'} \frac{\partial A}{\partial t_1} \right] e^{i\theta}

- \frac{E^{(2)}_0 (1 + \sigma_2^2)}{2 \sigma_2} \left( \Lambda - \frac{k \sigma_1^2}{\sigma_2^2} \right) \frac{\sinh(2(ky - b'))}{\cosh 2b'} - \Lambda^2 e^{2i\theta} + c.c.,$$

(20)

where $\xi_2$ (which represents the induced mean motion or the zero frequency correction to slow modulation of the fundamental mode), and $B^{(j)}_V$ are real functions of the slow scales $x_1, x_2, t_1,$ and $t_2$ to be determined by considering the equations of higher orders, and
\[
\begin{align*}
&= \frac{2\omega}{k} \left(1 + \frac{\rho_1}{\sigma_1}\right) \left(\frac{\partial A}{\partial t_2} + v_g \frac{\partial A}{\partial x_2}\right) - i \left(1 + \frac{\rho_1}{\sigma_2}\right) \frac{\partial^2 A}{\partial t_2^2} - \frac{2i}{k} \left(1 + \frac{\rho_1}{\sigma_1}\right) + \omega \left(1 + \frac{\rho}{\sigma_1\sigma_2} - \frac{\rho_1}{S_2^2}\right) \\
&= \frac{\partial^2 A}{\partial x_2 \partial t_1} + i \left[\omega^2 \left(1 + \frac{\rho_1}{\sigma_2}\right) - \frac{\rho_1(b - b')}{S_2^2} \left(\frac{a'}{\sigma_1} - \frac{b'}{\sigma_2}\right)\right] - \frac{1}{k} \left[a'(1 - \rho + 3k^2) + \frac{\sigma_1}{4}(1 - \rho + 13k^2)\right] \\
&+ \frac{V_2^2}{2} \frac{\partial^2 A}{\partial x_2^2} + \frac{2\beta_1}{\sigma_2} \left(1 - \frac{\beta_1}{\sigma_2}\right) \left(2 - \frac{b^2}{S_2^2}\right) + 2\beta_1 \left(V_2^2 b' + V_1^2 a'\right) + 4 \left(V_2^2 + V_1^2\right) \frac{\partial^2 A}{\partial x_2^2} \\
&+ iG_3 A^2 A + i \left[\frac{G_2}{\sigma_1} + 2\omega \left(1 + \frac{\rho_1}{\sigma_1} + \frac{\partial B_1^{(2)}}{\partial x_2}\right)\right] A = 0,
\end{align*}
\]

where

\[
G_3 = \omega^2 \left\{ \lambda \left(\frac{\rho_1}{\sigma_1} + \frac{\rho_2}{\sigma_2} + \frac{1 - 2\sigma_2^2}{\sigma_1}\right) + k^2 \frac{\rho_1^2}{\sigma_1^2} \left(\frac{2 - 7\sigma_2^2}{\sigma_2^2} + \frac{1 - \sigma_2^2}{\sigma_1^2}\right)\right\}
\]

Furthermore, from the non-secularity condition for \(\eta_3\), we have

\[
\frac{\partial \xi_2}{\partial t_1} \left[1 + \frac{\rho_1}{\sigma_2}\right] \frac{\partial B_1^{(j)}}{\partial x_2} + 2\omega \left(1 - \frac{\sigma_1}{\sigma_2}\right) \frac{\partial A}{\partial x_2} = 0.
\]

If we assume that \(\xi_2\), \(B_1^{(j)}\) as well as \(A\) depend on \(x_1\) and \(t_1\) only through \(\xi = x_1 - v_g t_1\), (26) yields

\[
\left\{a'b'(1 - \rho) - kv_g^2 (b' + \rho a')\right\} \frac{\partial B_1^{(j)}}{\partial \xi} =
\]

\[
k \left\{2\omega \rho_1 \frac{1}{\sigma_1 + \sigma_2} \right\} v_g^2 (\rho, 1) - 2\omega \sigma_1 (1 - \rho)(b', -a'\sigma_1/\sigma_2) + v_g (-b', -a'G_2) |A|^2 + f_j(x_2, t_2),
\]

where we assume \(v_g^2 \neq ab(1 - \rho)(b + \rho a)\). The slow functions \(f_j(x_2, t_2)\) in (27) are to be determined under appropriate boundary and/or initial conditions of the problem under consideration. Hereafter, however, we omit these terms, since they can be eliminated from the final result by a simple transformation [38] and cause no effect on the stability characteristics. Then \(\xi_2, \frac{\partial B_1^{(j)}}{\partial x_2}\) in (25) can be expressed in terms of \(A\). Using (23) and (24), introducing the expressions for \(\xi_2, \frac{\partial B_1^{(j)}}{\partial x_2}\) into (25), and assuming that \(A\) depends on \(x_2\) and \(t_2\) through \(v_g^2 = v_2^2 = v_g t_2\) and \(\tau = t_2\), we obtain finally a nonlinear Schrödinger equation

\[
i \frac{\partial A}{\partial \tau} + \frac{\partial^2 A}{\partial \xi^2} + v |A|^2 A = 0,
\]

where

\[
\mu = \frac{1}{2} \frac{\partial v_g}{\partial k}
\]
where $v$ is the nonlinear intercation coefficient; it should be noted that this coefficient becomes infinite and the perturbation scheme becomes invalid for values of $k$, $p$, $a$ and $b$ which satisfy $v = ab(1-p)/(b+pa) = 0$, which indicates that the group velocity of the wave train (i.e. $v_g$) coincides with the phase velocity of the infinitely long waves (i.e. $\sqrt{ab(1-p)/(b+pa)}$). In this case we have to modify the perturbation expansion so as to avoid the trouble of unboundedness. It is interesting to note that this case is expected to indicate a kind of resonant interaction between the group and phase velocities (i.e. between the short and the long waves); such a modification may be possible following the same lines as in [39].

### 6. Numerical Discussion

Equation (28) describes the nonlinear self-modulation of the capillary-gravity waves on liquid layers of uniform depths. It is interesting to note that the two coefficients $\mu$ and $\nu$ are responsible for the modulational instability of a nonlinear plane wave solution of (28). The original wave train is stable or unstable if $\mu \nu > 0$ or $\mu \nu < 0$. The stability characteristics change critically depending on the values of $\rho$, $a$, $b$, and $k$. The stability chart in the $k$-$p$ plane is divided into stable and unstable regions bounded by the curves.

(a) In the case of $\mu \nu > 0$

$$A_B(\zeta, \tau) = A_0 \text{sech}(K_+ \zeta) \exp\left(i(vA_0^2/2)\tau\right)$$

with

$$K_+ = \sqrt{vA_0^2 / 2 \mu}.$$

This convex envelope wave is called a bright soliton.

(b) In the case of $\mu \nu < 0$

$$|A_D(\zeta, \tau)|^2 = A_\infty^2 - A_0^2 \text{sech}^2(K_- \zeta)$$

with

$$K_- = -vA_0^2 / 2 \mu, \quad A_\infty > A_0 > 0.$$

This concave envelope wave is called a dark soliton or envelope hole. In this case, a shock type solution called a phase jump also exists:

$$A_P(\zeta, \tau) = A_0 \tanh(K_- \zeta) \exp\left(i(A_0^2 \nu)\tau\right).$$

The above envelope waves all stand steadily in $(\zeta, \tau)$ space.

If in the nonlinear Schrödinger equation (28), we take the limit when $ka \to -\infty$ and $kb \to \infty$ (i.e. $\sigma_1$, $\sigma_2 \to 1$), we recover the results obtained earlier by Mohamed and Elshehawey [29]. The results of Qi-su [7] for the corresponding pure hydrodynamical case can be obtained by setting $E_0^{(i)} = 0$ or $E_0^{(l)} = E_0^{(s)}$ in the nonlinear Schrödinger equation (28). We should also remark here that for ideal fluids, in the limit of no capillarity, (28) recovers the result for the gravity waves obtained earlier by Tanaka [8].

As we mentioned before, the original wave train is stable or unstable if $\mu \nu > 0$ or $\mu \nu < 0$. The stability chart in the $k$-$p$ plane is divided into stable and unstable regions bounded by the curves.

$$\mu = 0$$

and

$$\nu = 0.$$  

We observe from (35) and (21) that $v$ changes sign across the transition curves

$$G_1 = 0$$

and

$$kv_0^2(b' + pa') - a'b'(1 - \rho) = 0,$$

which represent the third and fourth transition curves in the stability diagrams. We note from (29) and (30) that the stability of the system does not depend on which one of the fluids has a larger dielectric constant. We computed the relations (34)-(37) for different values of $V_0^2$ in the three cases $a' = b'$, $a' > b'$ and $a' < b'$, respectively, for values of the density ratio $\rho \leq 1$. In the graphs, we plot the neutral stability conditions (34)-(37). Figures 3–9 represent the stability diagrams in the $k$–$\rho$
plane, due to the nonlinearity effect and the presence of the electric field, drawn for the three cases \( a' = b' \), \( a' > b' \) and \( a' < b' \), respectively. Figures 3–5 represent the case \( a' = b' \) (i.e. with two equal depths, where \( a' = b' = 0.9 \)). In this case the resulting curves from (34)–(37) are represented by the solid, dotted, dashed and dot-dashed curves, respectively, and we notice that there are two dotted curves which correspond to (35), we refer to them as the upper and lower parts of the dotted curve.

In Fig. 3a, where \( V_1^2 = V_2^2 = 0 \) (the pure hydrodynamical case), there are three stable regions between the curves, the first region \( S_1 \) is above the upper part of the second curve, while the second stable region \( S_2 \) is between the first and the fourth curves, and the third region \( S_3 \) is between the third curve and the lower part of the second curve, respectively. There are three unstable regions too, the first region \( U_1 \) is between the upper part of the second curve and the fourth curve, the second unstable region \( U_2 \) is between the first and the third curves, and the third region \( U_3 \) is under the lower part of the second curve. In Figs. 3b, c, where \( V_1^2 = 0.009 \) and \( V_2^2 = 0.009, 0.04 \), respectively, we note that the solid curve goes up slightly, creating a new unstable region \( U_4 \) which increases due to increase the electric field values, and the regions \( S_1, S_3 \) decrease, while region \( S_2 \) increases; and the region \( U_1 \) decreases while the regions \( U_2 \) and \( U_3 \) increase. We note also that the upper part of the dotted curve and the fourth curve coincide at \( p \geq 0.8 \) and also the lower part of the dotted curve and the second curve.

In Figs. 4 we have \( V_1^2 = 0.009 \) and \( V_2^2 = 0.08, 0.15 \) and 0.3, respectively. The resulting curves here have the same behaviour as the curves in Fig. 3, but in addition we find that region \( U_4 \) increases and region \( S_2 \) in Fig. 4a is split into two regions \( S_3 \) as in Figs. 4b, c, where the first region decreases and the second one increases due to the increase of the electric field values. A new stable region \( S_4 \) appears in Figure 4c.

In Fig. 5a, where \( V_1^2 = 0.009 \) and \( V_2^2 = 0.5 \), the curves still have the same behaviour as in Fig. 4, and we notice that the new stable region \( S_4 \) increases due to the increase of the electric field. In Figs. 5b, c, where \( V_1^2 = 0.009 \) and \( V_2^2 = 0.8, 1.0 \), respectively, we find that the solid curve drops and changes the situation mentioned in the previous figures where regions \( S_4, U_4 \) disappear and regions \( S_1 \) and \( S_2 \) increase. The first region of \( S_3 \) diminishes and the second region increases, and region \( U_2 \) changes its place. A new unstable region \( U_5 \) appears due to the increase of the electric field values. In Fig. 5c, we find that regions \( S_1, U_2 \), and \( U_5 \) increase, while regions \( U_1, U_3 \), and \( S_3 \) decrease, creating a new stable region \( S_5 \).

Figures 6 and 7 represent the case \( a' > b' \) (i.e. when the lower depth is larger than the upper one, where \( a' = 0.9, b' = 0.3 \)). In Fig. 6a where \( V_1^2 = V_2^2 = 0 \), we notice that we get curves quite similar to those in Fig. 3a for the case \( a' = b' \), but in this case we have four stable regions and four unstable regions due to the intersection of the upper part of the second curve with the fourth curve (at \( \rho = 0.12 \)) producing two unstable regions \( U_1 \) instead of the first unstable region \( U_1 \) in the previous case, and
we call them the first and second parts of the first unstable region $U_1$. Also the intersection of the lower part of the second curve with the third curve (at $\rho = 0.13$) produces two stable regions $S_3$ instead of the third stable region, and we call them too the first and second parts of the third stable region $S_3$.

In Figs. 6b, c, where $V_1^2 = 0.009$ and $V_2^2 = 0.005$ and 0.009, respectively, we notice that as in the previous case, the increase of the electric field slightly decreases the first stability region $S_1$ and the first part of the third region $S_3$, while it increases the second stability regon $S_2$ and the second part of the third region $S_3$. The two parts of the first unstable region $U_1$ decrease and the second unstable region $U_2$ increases while the third unstable region $U_3$ decreases. The solid curve turns out at $\rho = 0.9$, creating a new unstable and stable regions $U_4$ and $S_4$, respectively, under it with further increasing of the electric field values as shown in Figures 7.

In Figures 7, where $V_1^2 = 0.009$ and $V_2^2 = 0.03, 0.055$ and 0.08, respectively, we note that the behaviour of the
Fig. 6. Stability diagram for the nonlinearity effect in the $k$-$\rho$ plane for $\rho \leq 1$ and $a' = 0.9$, $b' = 0.3$. The solid, dotted, dashed, and dot-dashed curves represent respectively equations (34)-(37), when (a) $V_1^2 = V_2^2 = 0$, (b) $V_1^2 = 0.009$, $V_2^2 = 0.005$, and (c) $V_1^2 = 0.009$, $V_2^2 = 0.009$.

different curves and regions is still as before except that the solid curve goes up slightly, creating two new regions, one being unstable $U_4$ and the other one stable $S_4$ under it as in Figure 7a. It is clear from Figs. 7b, c that the increase of the electric field increases the new unstable region $U_4$ while it decreases the new stable region $S_4$; i.e. the normal electric field has a destabilizing effect in the two new regions.

Figures 8 and 9 represent the case $a' < b'$ (i.e. when the lower depth is smaller than the upper one, where $a' = 0.3$, $b' = 0.9$). Figure 8 are drawn for the electric field $V_1^2 = 0.009$ and $V_2^2 = 0.005, 0.02$ and $0.06$, respectively. We notice that the curves have the same behaviour as in the case $a' = b'$ (i.e. we have three stable regions $S_1$, $S_2$, $S_3$ and three unstable regions $U_1$, $U_2$, $U_3$), and due to the increasing of the electric field as in Fig. 8b, the first and the third stability regions $S_1$ and $S_3$ decrease while the second stability region $S_2$ increases; and the three instability regions $U_1$, $U_2$ and $U_3$ increase. Also the solid curve goes up slightly, creating two new regions, one is
Fig. 8. Stability diagram for the nonlinearity effect in the $k$-$p$ plane for $p < l$ and $a' = 0.3, b' = 0.9$. The solid, dotted, dashed, and dot-dashed curves represent respectively equations (34)–(37), when $V_f = 0.005$ and (a) $V_2 = 0.009$, (b) $V_2 = 0.02$, and (c) $V_2 = 0.06$.

Fig. 9. Stability diagram for the system considered in Fig. 8, when $V_f = 0.009$, but with (a) $V_2 = 0.1$, (b) $V_2 = 0.2$, and (c) $V_2 = 0.3$.

We find that in addition to the previous behaviour of the curves, the three new regions $S_5$, $S_6$, and $U_5$ increase due to increasing the electric field, and also a new stable region $S_7$ appears.

Thus the effect of the electric field is the same in all the cases for small values of the electric field; and for a fixed value of $V_f$ and after a definite value of $V_2$, the effect of the electric field is different from one case to another, and it is stronger in the case when $a' > b'$. 
7. Amplitude Modulation of Standing Waves

When \( k \) tends to \( k_c \), \( \omega \) tends to zero, so that the group velocity \( v_g \), and the coefficients \( \mu \) and \( \nu \) given by (29) and (30), respectively, become infinite. Therefore the preceding obtained results are no longer valid near the cut-off wavenumber \( k = k_c \), and we need to modify the preceding analysis to obtain the suitable equation for the complex amplitude near the cut-off wavenumber. Since we are concerned with waves near \( k = k_c \) and \( \omega = 0 \), the carrier wave is not travelling, but standing, hence we choose the starting solutions to the first order problem as

\[
\eta_1 = i \sigma_1 A e^{ik_c x_0} + \text{c.c.,} \quad (38)
\]

\[
\phi_1^{(j)} = B_1^{(j)} , \quad j = 1, 2 , \quad (39)
\]

\[
\psi_1^{(1)} = i E_0^{(1)} \frac{\sinh (k_c y + a')}{\cosh a'} e^{ik_c x_0} + \text{c.c.,} \quad (40)
\]

\[
\psi_1^{(2)} = - i E_0^{(2)} \frac{\sigma_1 \sinh (k_c y - b')}{\cosh b'} \frac{\partial A}{\partial t_1} e^{ik_c x_0} + \text{c.c.} \quad (41)
\]

instead of (9)–(13), where \( a' = k_c a \) and \( b' = k_c b \).

Proceeding as before, we can get the uniformly valid solution of the second order problem in the form

\[
\eta_2 = \left( A A^2 e^{2ik_c x_0} + \text{c.c.} + \varepsilon_2 \right) , \quad (42)
\]

\[
\phi_2^{(1)} = \frac{i \cosh (k_c y + a')}{k_c \cosh a'} \frac{\partial A}{\partial t_1} e^{ik_c x_0} + \text{c.c.} + B_2^{(1)} , \quad (43)
\]

\[
\phi_2^{(2)} = - \frac{i \cosh (k_c y - b')}{k_c \cosh b'} \frac{\sigma_1 A}{\sigma_2} \frac{\partial A}{\partial t_1} e^{ik_c x_0} + \text{c.c.} + B_2^{(2)} , \quad (44)
\]

\[
\psi_2^{(1)} = \frac{E_0^{(1)} (1 + \sigma_1^2)}{2 \sigma_2} \left( A + k_c \sigma_1^2 \right) \frac{\sinh 2(k_c y + a')}{\cosh 2a'} A e^{2ik_c x_0} + \text{c.c.} , \quad (45)
\]

\[
\psi_2^{(2)} = - \frac{E_0^{(2)} (1 + \sigma_1^2)}{2 \sigma_2} \left( A - k_c \sigma_2^2 \right) \frac{\sinh 2(k_c y - b')}{\cosh 2b'} A e^{2ik_c x_0} + \text{c.c.} , \quad (46)
\]

where

\[
A = \frac{k^2 \sigma_2^2}{2} \left[ \frac{V_2^2 \left( \frac{2}{\sigma_2^2} + 1 \right)}{S_2^2} - V_1^2 \left( \frac{2}{\sigma_1^2} + 1 \right) \right] \cdot \left( V_2^2 \sigma_2^2 + V_1^2 \sigma_1^2 - 3k_c \right)^{-1} \quad (47)
\]

under the non-secular condition

\[
\frac{\partial A}{\partial x_1} = 0 , \quad (48)
\]

which shows that the complex amplitude \( A \) is independent of the first slow scales \( x_1 \). This should be compared with (23) for the travelling waves.

Let us now proceed to the third order problem. After straightforward calculations, the requirement that the third order perturbation be non-secular yields

\[
i \frac{\partial A}{\partial x_2} + \mu^* \frac{\partial^2 A}{\partial t_1^2} = \nu^* | A |^2 A + RA , \quad (49)
\]

where

\[
\mu^* = - \frac{1}{2} \frac{\partial^2 k}{\partial \omega^2} \bigg|_{\omega=0} \quad , \quad (50)
\]

\[
\nu^* = \frac{k_c}{2} \frac{\partial^2 k}{\partial \omega^2} \bigg|_{\omega=0} \left( \frac{1}{\sigma_1} + \frac{\rho}{\sigma_2} \right) \quad , \quad (51)
\]

\[
R = - \frac{k^2}{2} \frac{\partial^2 k}{\partial \omega^2} \bigg|_{\omega=0} \left( \frac{1}{\sigma_1} + \frac{\rho}{\sigma_2} \right) \quad , \quad (52)
\]

where \( E(x_2, t_2) \) is a constant in the slow scales \( x_1 \) and \( t_1 \).

Equation (49) is also a nonlinear Schrödinger equation, with time and space variables interchanged and hence furnishes the amplitude modulation of standing waves. The linear interaction term \( RA \) in (49) causes on-
ly a phase shift and thus can be removed by using the transformation

\[ A(x_2, t) = A \exp \left( -i \int R(x'_2) \, dx'_2 \right). \]  

(53)

The nonlinear Schrödinger equation (49) now takes the form

\[ i \frac{\partial A}{\partial x_2} + \frac{\mu}{2} \frac{\partial^2 A}{\partial t^2} = v^* |A|^2 A. \]  

(54)

We now examine the plane wave solution of the form

\[ A(x_2, t) = B_0 e^{i(Kx_2 - \Gamma t)} e^{-\int R(x'_2) \, dx'_2}, \]  

(55)

where \( B_0 \) is a complex constant, and \( K \) and \( \Gamma \) are real constants representing respectively the wavenumber and frequency shifts.

On substituting (55) into (54), we get the dispersion relation

\[ \Gamma^2 = -\left( K + v^* |B_0|^2 \right) / \mu^*. \]  

(56)

For \( \Gamma \) to become imaginary, we must have \( K < v^* |B_0|^2 \). Combining (55) with the carrier wave, we can determine the "nonlinear" cut-off wavenumber as

\[ k_n = k_c \left[ 1 + \frac{\varepsilon |B_0|^2}{2} \frac{\partial^2 k}{\partial \omega^2} \bigg|_{\omega=0} \left( \frac{1}{\sigma_1} + \frac{\rho}{\sigma_2} \right)^{-1} \right] \]

\[ \cdot \left[ \frac{3}{2} k^4 \sigma_1^2 - k^2 k_c \left( \frac{V_2^2}{\sigma_2^2} + \frac{1}{S_2^2} \right) - V_1^2 \left( \frac{2}{\sigma_1^2} + \frac{1}{S_1^2} \right) \right] \]

\[ + k^2 \sigma_1^2 \left[ V_2^2 \left( \frac{1}{S_2^2} - 1 \right) - V_1^2 \left( \frac{1}{S_1^2} - 1 \right) \right] \]

\[ - \frac{\varepsilon}{x_2} \int R(x'_2) \, dx'_2, \]  

(57)

which shows that the nonlinear cut-off wavenumber \( k_n \) depends sensitively on the initial condition with respect to \( t_1 \). It is clear that the bandwidth of the spectrum is of \( O(\varepsilon^4) \) in the frequency space for the standing wave, i.e. it is of \( O(\varepsilon^4) \) in the wavenumber space.

8. Conclusions

In this work, we have investigated the nonlinear electrohydrodynamic stability of interfacial capillary-gravity waves of two superposed dielectric fluids with uni-
pear, one of them is stable and the other is unstable, and they increase due to increase the electric field in the case $a' = b'$, while the new stable region decreases and the unstable region increased if $a' < b'$. For a fixed value of $V_f^2$ and after a definite value of $V_2^2$, the distribution of stable and unstable regions in the first case $a' = b'$ changes and two more stable and unstable regions appear and increase; while in the third case $a' < b'$, only two more stable and unstable regions appear and increase due to increase the electric field.

In the second case $a' > b'$, the stability diagram is divided by four stable and four unstable regions, and the effect of the electric field is destabilizing in the first stable and the first part of the third stable region, and stabilizing in the second stable and the second part of the third stable region. Two newly regions appear, one of them is unstable and the other one is stable, and the electric field increases the first new region and decreases the second new one.

Thus the effect of the electric field is the same in all the cases for small values of the electric field; and for a fixed value of $V_f^2$ and after a definite value of $V_2^2$, the effect of the electric field is different from one case to another, and it is stronger if $a' > b'$. Therefore the normal electric field, in the presence of surface charges at the interface, creates some new regions of stability and instability, which were hidden in the corresponding case when there are no charges at the interface between the two fluids [34].

We also recovered some previous work and limiting cases corresponding to, e.g. the pure hydrodynamical case, the case of two semi-infinite fluids, the case of capillary waves, etc.; and the results of the linear theory are confirmed. We treated our problem too in the case of standing waves near the cut-off wavenumber, which is governed also by a nonlinear Schrödinger equation with the roles of time and space are interchanged. We finally determined the nonlinear dispersion relation and the nonlinear effect on the cut-off wavenumber.

**Appendix**

The solutions of the third order problem are given by

$$
\eta_3 = \left[ \frac{i(4a' + \sigma_1)}{4k^2} \frac{\partial^2 A}{\partial x_1^2} + \frac{i(a' + \sigma_1)}{k} \left\{ \frac{1}{\omega} \frac{\partial^2 A}{\partial x_1 \partial t_1} + i \frac{\partial A}{\partial t_1} \right\} \frac{-i\sigma_1}{\omega} \frac{1}{\omega} \left\{ \frac{1}{\omega} \frac{\partial^2 A}{\partial t_1^2} + i \frac{\partial A}{\partial t_1} \right\} \right]
$$

(A.1)

$$
\phi_3^{(1)} = -\frac{\omega}{k \cosh a'} \left[ \frac{y^2}{2} + ay + \frac{1}{4k^2} \right] \cosh(ky + a') \frac{\partial^2 A}{\partial x_1^2} + i(y + a) \sinh(ky + a') \frac{\partial A}{\partial x_2} \right] e^{i\theta} + NSPT + c.c. + \xi_3,
$$

(A.2)

$$
\phi_3^{(2)} = \left[ \frac{\omega}{k \cosh b'} \frac{\sigma_1}{\sigma_2} \right] \left\{ \frac{y^2}{2} + by + \frac{1}{4k^2} \right\} \cosh(ky + b') \frac{\partial A}{\partial x_2} \right] + \frac{\omega \psi_1}{k^2 \sigma_2^2} \frac{a' - b'}{\sigma_1 - \sigma_2} \frac{\partial^2 A}{\partial x_1^2} + i \frac{\omega \psi_1}{k^2 \sigma_2} \frac{a' - b'}{\sigma_1 - \sigma_2} \frac{\partial A}{\partial x_2} + \frac{\omega(\sigma_1 + \sigma_2)}{\sigma_2^2}
$$

(A.3)
\[ \psi_3^{(1)} = -iE_0^{(1)} \left[ \frac{\sinh(ky + a')}{\cosh a'} \left( \int k \frac{1}{\omega} \frac{\partial^2 A}{\partial x_1 \partial t_1} \right) + \int \frac{1}{\omega} \frac{\partial^2 A}{\partial x_2 \partial t_2} \right] - \frac{k}{\omega} \frac{\partial B_1^{(1)}}{\partial x_1} A \right] \]

\[ + \left[ \frac{1}{k} \frac{\partial^2 A}{\partial x_1^2} + \frac{1}{\omega} \frac{\partial^2 A}{\partial x_1 \partial t_1} + i \frac{\partial A}{\partial x_2} \right] \left( \frac{y + a}{\cosh (ky + a')} \cosh a' \right) \]

\[ + \left( \frac{y^2}{2} + ay + \frac{1}{4k^2} \right) \sinh(ky + a') \frac{\partial^2 A}{\partial x_1^2} e^{i\theta} + NSPT + c.c., \]  
(A.4)

\[ \psi_3^{(2)} = iE_0^{(2)} \sigma_1 \left[ \frac{\sinh(ky - b')}{\cosh b'} \left( \int \frac{a' - b'}{\sigma_1} - \frac{1}{\sigma_2} \frac{\partial^2 A}{\partial x_1 \partial t_1} + \int \frac{1}{\omega} \frac{\partial^2 A}{\partial x_2 \partial t_2} + i \frac{\partial A}{\partial x_2} \right) + \frac{k}{\sigma_1} \right] \]

\[ \left( \int \frac{2 - \sigma_2}{\sigma_1} - \frac{1}{\sigma_1} \frac{\sigma_2}{\sigma_2} \right) \right] A^2 \Delta - k \left[ \frac{\sigma_1 + \sigma_2}{\sigma_1 \sigma_2} \right] \frac{\partial B_3^{(1)}}{\partial x_1} A \right] \]

\[ (A.5) \]

\[ + \left[ \frac{1}{k} \frac{a' - b'}{\sigma_1} + 1 \right] \frac{\partial^2 A}{\partial x_1^2} + \frac{1}{\omega} \frac{\partial^2 A}{\partial x_1 \partial t_1} + i \frac{\partial A}{\partial x_2} \left( \frac{y - b}{\cosh (ky - b')} \cosh b' \right) \]

\[ + \left( \frac{y^2}{2} - by + \frac{1}{4k^2} \right) \sinh(ky - b') \frac{\partial^2 A}{\cosh b'} \frac{\partial A}{\cosh b'} \] 

\[ e^{i\theta} + NSPT + c.c., \]

where here and hereafter NSPT represents terms that do not produce secular terms, and \( \xi_3, B_3^{(1)} \) and \( B_3^{(2)} \) are real functions of the slow scales \( x_1, x_2, t_1 \) and \( t_2 \) to be determined by considering the equations of higher orders.

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