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The non-linear breakup of an inviscid liquid jet

D.P. BUSKER

Volvo Car,Born, The Netherlands

and

A.P.G.G. LAMERS

Department of Mechanical Engineering, Eindhoven University of Technology, Eindhoven, The Netherlands

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Abstract. A liquid jet originating from a nozzle with radius \( r_0^* \) breaks up into droplets in consequence of disturbances of certain frequencies, depending on the fluid properties and the nozzle geometry. A theoretical model is developed to describe the growth of these disturbances at the jet surface. The model is based on the inviscid and irrotational flow governed by the Laplace equation together with the kinematical and dynamical conditions at the free surface of the jet. A comparison is made between the model and experimental data from literature. The model predicts a dependence on the disturbance amplitude of the breakoff mode. Contrary to other experimental results, the model predicts satellites (i.e. smaller droplets between the main larger ones) at wavelengths exceeding a critical value of \( \frac{1}{2} \times 2\pi r_0^* \). The disturbances grow at wavelengths more than the theoretical bound of \( 2\pi r_0^* \). Discrepancies with experimental data are possible because of the neglect of the effect of viscosity in the theory. It is shown that the effect of viscosity on the jet can be neglected under certain conditions.

1. Introduction

A liquid jet originating from a nozzle is sensitive to disturbances. Disturbances of certain frequencies cause the jet to break up into a series of successive droplets. There are two theoretical methods to investigate the behaviour of a disturbed liquid jet. The "spatial instability" method describes the disturbance of the jet surface as a travelling wave in axial direction. The "temporal instability" method describes the surface disturbance of the jet as a standing wave on an infinitely long cylinder with the nozzle at infinity.

The model presented in this paper is based on spatial instability and describes the jet form close to the nozzle in the form of travelling waves with harmonic influences. The Laplace equation together with the dynamical and kinematical boundary condition for the free surface are used to describe the radius of the jet and the velocity. The model is mathematically simplified by neglecting both the effects of viscosity and the surroundings. It is possible to approximate the solution of the equations for the radius of the jet by the solution of a simplified fluid dynamical theory, the Cosserat theory, which is a simplified one-dimensional theory. In this article we present an approximation for radius and velocity by Taylor series expansions with respect to the disturbance amplitude.
The breakup process depends on surface tension $\sigma^*$, density $\rho^*$ and nozzle radius $r_0^*$ and initial disturbance amplitude $\delta_0^*$. (Dimensional variables are indicated with an asterisk throughout.) The disturbances grow with time and distance from the nozzle. Piezo crystals or mechanical vibrators can be used as sources when applied to the jet surface, the velocity or the pressure distribution in the jet.

Previously published theoretical models are of a more limited use as the model presented in this article, either because the stability analysis ignores higher harmonic effects or because these models are based on a simplified fluid dynamical theory. The spatial instability method describes the physical reality better than the temporal instability method does, whereas the first method does not impose periodic axial demands on the jet. The model based on spatial instability shows that satellites can break up before or after a main drop, dependent on the disturbance amplitude. A condition is derived by which the influences of the viscosity can be estimated and eventually neglected.

2. Literature

The first mathematical model was published by Rayleigh (1878). This model was a stability analysis for infinitely small disturbances of an inviscid jet based on the temporal instability method. Weber (1931) and recently Sterling and Sleicher (1975) extended this analysis to an aerodynamically influenced viscous jet. According to these linearized models the only dependent variable for the breakup process is the breakup time.

Approximately 1960 these linearized theories were found inadequate to describe the phenomenon accurately. Only the breakup time could be reasonably well compared with experiments. At present the formation of satellites is considered to be more characteristic for the phenomenon than the breakup length or breakup time. Yuen (1968) and Lafrance (1974, 1975) have obtained analytical approximations for an inviscid jet ignoring influences of the environment on the basis of temporal instability and mass conservation for higher harmonic disturbances. They arrived at establishing the formation and the existence of satellites. Rutland and Jameson (1970) compared experimental data with numerical results based on Yuen's model. By coupling the initial and final volume of the droplets they succeeded in calculating diameters of representative satellites and main drops. They concluded that in case of wavelengths $\lambda^* \leq \frac{10}{3} \times 2\pi r_0^*$ no satellites could be formed. With a modified model, Lafrance (1974) showed that no satellites were formed for $\lambda^* \leq \frac{10}{8} \times 2\pi r_0^*$. These conclusions are certainly not correct in the case of viscous jets. Chaudhary and Redekopp (1980) and Chaudhary and Maxworthy (1980a, 1980b) used a comparable model and performed experiments on the jet behaviour and satellite drop formation.

In 1973 Keller et al. used spatial instability in combination with a stability analysis to describe the breakup process of a jet. As a result, the description of the breakup process is improved. From a mathematical viewpoint the disturbances behave like waves progressing on the jet surface. Bogy (1978, 1979a, 1979b) expanded the Cosserat theory of Green (1976) to develop a new analytical model for higher harmonics based on spatial instability. The Cosserat theory is a one-dimensional theory simplifying the flow in the jet, especially in the radial direction.

The model presented in this paper is an extension of the models referred to above. It describes the break-up process of a jet with the spatial instability method and mass conservation including higher harmonic effects. The jet form is approximated by a Taylor series expansion with respect to the disturbance amplitude. The geometry of the jet, the velocity components in it, the breakup length and the breakup mode of satellites before or behind a main drop are calculated with our model. Computation of the last mentioned phenomenon is
not found in earlier published articles. The model still has physical limitations on account of
the neglect of viscosity and environmental effects such as gravity, mass or heat transfer.

3. Mathematical model

3.1. General description

A semi-infinite axisymmetric jet with liquid density $\rho^*$ and surface tension $\sigma^*$ emerges from
a nozzle of radius $r_0^*$. The uniform velocity $v_0^*$ at the nozzle is harmonically disturbed. As a
result the radius $R^*$ of the jet is a function of the axial coordinate $z^*$ and time $t^*$. Fig. 1 shows
the geometry of the jet. The characteristic length and time are $r_0^*$ and $t_0^* = r_0^*/v_0^*$ respectively.
The breakup process of the jet has a characteristic timescale $t_1^* = (\rho^* r_0^* / \sigma^*)^{1/2}$. The charac-
teristic dimensionless number is the Weber number which is defined as the quadratic ratio of
the characteristic time scales $t_1^*$ and $t_0^*$:

$$\text{We} = \left( \frac{t_1^*}{t_0^*} \right)^2 = \rho^* r_0^* (v_0^*)^2 / \sigma^*.$$  \hspace{1cm} (1)

The influence of surroundings, viscosity and gravity are neglected. Rotational symmetry and
irrotational flow in the jet are assumed. The flow in the jet is described by means of the Laplace
equation, a pressure continuity condition at the jet boundary and a containment boundary
condition, in the sense that all liquid remains within the jet surface. The nozzle diameter is
constant and the axial jet velocity at the nozzle is sinusoidally disturbed.

The local jet radius $R^*$ and the velocity potential $\Phi^*$ are approximated by series expan-
sions. The non-linear differential equations are reduced to a set of linear differential equations;
these are solved up to the third order to be able to describe satellites. From the homogeneous
differential equations the characteristic dispersion relation is derived. The dispersion relation is
solved numerically on a mainframe Burroughs. The breakup length $z_B^*$ and the local jet radius
$R^*(z^*, t^*)$ are calculated by the computer ($z^*$ is the axial distance). These results are
compared with the experimental data taken from literature (Chaudhary and Maxworthy, 1980a,
1980b).

3.2. Formulation

A velocity potential $\Phi^*$ with $u^* = \partial \Phi^*/\partial z^*$ and $v^* = \partial \Phi^*/\partial r^*$ is introduced. All varia-
bles are made dimensionless with respect to the characteristic time $t_0^* = r_0^*/v_0^*$ and the
characteristic length $r_0^*$. The Laplace equation in terms of dimensionless $\Phi$ is

$$\nabla^2 \Phi = \Phi_{rr} + r^{-1} \Phi_r + \Phi_{zz} = 0 \quad (0 \leq r \leq R, \ 0 \leq z \leq z_B),$$  \hspace{1cm} (2)

Fig. 1. Sketch of the geometry of a liquid jet emerging from a nozzle at $z^* = 0$. The entrance velocity is disturbed
sinusoidally. The local jet radius $R^*$ is a function of the axial coordinate and time. Under certain conditions the jet
breaks up into main drops, spaced at wavelength $\lambda^*$, and satellites at the breakup length $z_B^*$. 

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with \( z_B \) as the breakup length of the jet, and subscripts \( r \) and \( z \) indicating derivatives w.r.t. \( r \) and \( z \) respectively. The kinematical boundary condition for free surface of the jet is

\[
\Phi_r = R_z \Phi_z (r = R),
\]

where subscript \( t \) indicates a time derivative. The dynamical boundary condition is given by Busker (1983)

\[
\Phi_t + \frac{1}{2} \left( \Phi_r^2 + \Phi_z^2 \right) + \frac{1}{\text{We}} \left( \frac{1}{R (1 + R_z^2)^{1/2}} - \frac{R_z}{(1 + R_z^2)^{3/2}} \right) = \frac{1}{2} + \frac{1}{\text{We}} (r = R).
\]

The nozzle diameter is constant,

\[
R = 1 \quad (z = 0)
\]

and the jet velocity is disturbed sinusoidally according to

\[
\Phi_z = 1 + \delta_0 \cos \omega t \quad (z = 0, 0 \leq r \leq 1).
\]

Solutions of \( R \) and \( \phi \) are considered in the interval \( 0 \leq z \leq z_B \).

After the breakup point the jet is broken up into droplets and is no longer a continuous medium. Chaudhary and Redekopp (1980) started from the same equations ((2), (3) and (4)). Our approach, however, is to seek for solutions of the form

\[
R = 1 + R' = 1 + \delta_0 \eta_1 + \delta_0^2 \eta_2 + \delta_0^3 \eta_3 + O(\delta_0^4)
\]

and

\[
\phi(r, z, t) = z + \delta_0 \phi_1 + \delta_0^2 \phi_2 + \delta_0^3 \phi_3 + O(\delta_0^4).
\]

Here \( \eta_i \) is the \( i \)th order harmonic surface distortion, and \( \phi_i \) is the \( i \)th order harmonic velocity potential of the main disturbance with \( i = 1 \). The derivative of the velocity potential on the boundary is transferred to the undisturbed jet geometry \( R = 1 \) by a re-expansion of the velocity potential \( \phi \) into a series of the boundary transformation \( R' \). After substitution of the boundary transformation \( R' \) the terms of eqs. (2)–(6) can be expressed in terms of \( \delta_0, \delta_0^1, \delta_0^2 \) and \( \delta_0^3 \) respectively. After that, the equations are separated into powers of \( \delta_0 \). The zeroth-order solution gives the undisturbed cylindrical jet. The first-order equations are given by

\[
\nabla^2 \Phi_1 = 0 \quad (0 \leq r \leq 1),
\]

\[
\Phi_{1,r} - \eta_{1,r} - \eta_{1,z} = 0 \quad (r = 1),
\]

\[
- \Phi_{1,t} - \Phi_{1,z} + \frac{1}{\text{We}} (\eta_1 + \eta_{1,zz}) = 0 \quad (r = 1),
\]

with the following nozzle conditions

\[
\eta_1 = 0 \quad (z = 0),
\]

\[
\Phi_{1,z} = \cos \omega t \quad (z = 0, 0 \leq r \leq 1).
\]

From these first-order equations a stability rule, the dispersion relation, will be derived in section 3.3. The second-order (\( \delta_0^2 \)) and the third-order (\( \delta_0^3 \)) equations are given by Busker (1983).

3.3. Dispersion relation

Because of the first-order nozzle conditions (see (12), (13)) and eq. (9) which is defined at \( r = 0 \), a first-order solution is assumed which has the form

\[
\eta_1 = C \exp[i(\omega t - kz)] \quad (z, t \geq 0),
\]

\[
\Phi_1 = DI_0(kr) \exp[i(\omega t - kz)] \quad (z, t \geq 0, r \leq 1).
\]
Here, \( I_n(x) \) is the modified Bessel function of the first kind and \( n \)th order, \( \Phi_1 \) is a solution of the Laplace equation, and \( D \) is a constant. Substitution of \( \eta_1 \) and \( \Phi_1 \) into eqs. (10) and (11) produces the following characteristic equation:

\[
(\omega - k)^2 = \frac{k(k^2 - 1)}{\text{We}} \frac{I_1(k)}{I_0(k)}.
\]

(16)

This equation is called the dispersion relation, because it connects the frequency \( \omega \) to the wavenumber \( k \) at a given Weber number. The real part of the wavenumber is related to the wavelength \( \lambda \) by

\[
\text{Re}\{k\} = 2\pi/\lambda.
\]

(17)

The imaginary part of \( k \) determines the rate of growth of the disturbance. The dispersion relation has an infinite number of solutions \( k \) per \( \omega \). The solutions for \( k \) are functions of \( \omega \) and We. Keller et al. (1973) give the zeroth-order solutions for this equation

\[
(\omega - k)^2 = O(\text{We}^{-1}),
\]

(18)

leading to \( k = \omega + O(\text{We}^{-1/2}) \), and

\[
\frac{I_0(k)}{I_1(k)} = O(\text{We}^{-1}),
\]

(19)

leading to \( k = \pm i j_{0n} + O(\text{We}^{-1}) \) with \( I_0(\pm i j_{0n}) = 0 \). Keller et al. (1973) give an approximation for the solution of (18) and (19) respectively

\[
k_{1,2} = \omega \pm \left( \frac{\omega(\omega^2 - 1)}{\text{We}} \frac{I_1(\omega)}{I_0(\omega)} \right)^{1/2} + O(\text{We}^{-1}),
\]

(20)

\[
k_{3,4n} = \pm i j_{0n} + \frac{j_{0n}(j_{0n} j_{0n} + 1)}{\text{We}(\omega^2 + j_{0n}^2)} \left[ 2\omega j_{0n} \mp i(\omega^2 - j_{0n}^2) \right] + O(\text{We}^{-2}).
\]

(21)

Another solution is given by Boersma (1985)

\[
k_5 = \text{We} - (2\omega - \frac{1}{2}) + \frac{9}{8} + \omega - \frac{3\omega^2}{\text{We}} + O(\text{We}^{-2}).
\]

(22)

Solutions of eqs. (21) and (22) have wavelengths of the order of the Weber number which is large (\( \text{We} > 100 \)). Therefore they cannot occur in short jets. \( k_{1,2} \) will be the dominant breakup mode (Keller et al., 1973). For \( \omega > 0 \) only those solutions of \( k \), which have a positive phase-velocity \( c = \omega/\text{Re}(k) > 0 \) and a positive group velocity \( c_g = \partial \omega / \partial \text{Re}(k) > 0 \), are valid solutions because only then energy is transported downstream and the waves are travelling in the same direction (Bogy, 1978). The wavenumbers \( k_{1,2} \) come up to these conditions. Here it is assumed that the disturbance excitator is the only energy source in the system. The excitator is located at the nozzle and therefore energy goes downstream from the source in the same manner as the waves. Other assumptions are possible depending on the place of the disturbance source along the jet but they are beyond the scope of this article. A singular point exists for \( \omega = \omega_s \), because there is an abrupt changeover from complex to real wavenumber values. The dispersion relation has now two equal solutions. For \( \omega < \omega_s \) the two solutions are conjugate complex; see fig. 2. By approximation we find for the singular point \( \omega_s \)

\[
\omega_s \approx 1 - \frac{I_1(1)}{2 \text{We} I_0(1)} \approx 1 - \frac{0.221}{\text{We}}.
\]

(23)

Substitution of (23) in (20) gives for \( k_{1,2} \) the expression

\[
k_{1,2} \approx 1 + \frac{I_1(1)}{2 \text{We} I_0(1)} \approx 1 + \frac{0.221}{\text{We}}.
\]

(24)
Moreover the following relations hold
\[
\lim_{\omega \to \omega_c} \text{Im}\{k_{1,2}\} = 0, \quad \lim_{\omega \to \omega_c} \frac{\partial \text{Im}\{k_{1,2}\}}{\partial \omega} = \pm \infty, \quad \lim_{\omega \to \omega_c} \frac{\partial \text{Re}\{k_{1,2}\}}{\partial \omega} = 0.
\] (25)
The maximum growth rate occurs at \(\omega_{\text{opt}} = 0.4858^{1/2} \approx 0.6969\) for large Weber numbers.

3.4. Influence of the viscosity

The dispersion relation for a jet with viscosity \(\mu\) under the same conditions as an inviscid jet is given by Busker (1983)
\[
(\omega - k)^2 = + \frac{k(k^2 - 1)}{\text{We}} \frac{I_1(k)}{I_0(k)} - \frac{2i(\omega - k)k}{\text{Re}} \left(2k - \frac{I_1(k)}{I_0(k)}\right)
- \frac{4k^3}{\text{Re}^2} \left(k - \frac{I_1(l)I_1(k)}{I_1(l)I_0(k)}\right)
\] (26)
with \(l^2 = k^2 + i(\omega - k)\text{Re}\). In case of \(\omega < 1\) and by approximating \(I_1(k)/I_0(k) \approx \frac{1}{2}k + O(k^2)\) eq. (26) can be rewritten analogous to Weber (1931) as
\[
(\omega - k)^2 = \frac{k(k^2 - 1)}{\text{We}} \frac{I_1(k)}{I_0(k)} - \frac{3i(\omega - k)k^2}{\text{Re}}.
\] (27)
The solution for this equation is
\[
k_{1,2} = \omega \pm \left(\omega(\omega^2 - 1)\frac{I_1(\omega)}{I_0(\omega)}\right)^{1/2} \left(\frac{1}{\text{We}^{1/2}} + \frac{(3i\omega^2 + 2)}{(3i\omega^2 - 1)\text{Re}}\right).
\] (28)
The influence of the viscosity might be neglected if the influence of the viscous terms are less than 2% of the surface tension term, thus
\[
\left|\frac{(3i\omega^2 + 2)\text{We}^{1/2}}{(3i\omega^2 - 1)\text{Re}}\right| < 0.02.
\] (29)
For the examples of the experiment of Chaudhary and Maxworthy (1980a, 1980b) this gives in case of \( \omega_1 = 0.4313 \), \( \text{We} = 922.6 \) and \( \text{Re} = 1486 \) a value of \( \approx 3.7\% \) and in case of \( \omega_1 = 0.720 \), \( \text{We} = 330.9 \) and \( \text{Re} = 889.9 \) a value \( \approx 2.8\% \) (see section 4).

3.5. First-order solution

The dispersion relation provides the wavenumber values \( k_{1,2} \) for a given dimensionless frequency \( \omega_1 \) and Weber number which are used in the first-order solution (14) and (15).

\[
\eta_1 = P_1 \cos \Theta_1 + P_2 \cos \Theta_2, \quad (30)
\]

\[
\Phi_1 = Q_1 I_0(k_1r) \sin \Theta_1 + Q_2 I_0(k_2r) \sin \Theta_2, \quad (31)
\]

where \( \Theta_1 = \omega_1 t - k_1 z \), \( \Theta_2 = \omega_1 t - k_2 z \), and \( P_1 \), \( P_2 \), \( Q_1 \) and \( Q_2 \) are constant. Substitution into the kinematical boundary condition and the nozzle conditions produces the following first-order solution

\[
\eta_1 = \frac{P}{V} (\cos \Theta_1 - \cos \Theta_2), \quad (32)
\]

\[
\Phi_1 = \frac{P}{V} \left( \frac{(\omega_1 - k_1) I_0(k_1r)}{k_1 I_1(k_1)} \sin \Theta_1 + \frac{(\omega_1 - k_2) I_0(k_2r)}{k_2 I_1(k_2)} \sin \Theta_2 \right) \quad (0 \leq r \leq R = 1). \quad (33)
\]

Here \( P = I_1(k_1) I_1(k_2) \) and \( V = (\omega_1 - k_1) I_1(k_2) - (\omega_1 - k_2) I_1(k_1) \). Dependent of the choice of \( \omega_1 \), the values of \( k_1 \) and \( k_2 \) are real or complex conjugated (see fig. 2). Because of eqs. (32) and (33) the solutions of \( \eta_1 \) and \( \Phi_1 \) are always real.

3.6. Second- and third-order solutions

The second- and third-order solutions are derived from the second- and third-order differential equations together with the nozzle conditions. They consist of a particular part and a homogeneous part. The particular parts of the second-order boundary differential equations are composed of terms with \( \eta_1 \) and \( \Phi_1 \). In case of the third-order it is composed of terms with \( \eta_1 \), \( \Phi_1 \), \( \eta_2 \) and \( \Phi_2 \). By substituting the lower-order solutions in the particular part of the second- and the third-order differential equations a general form of the particular second- and third-order solution is derived. These general forms for the second- and third-order solutions are substituted in the homogeneous parts of the boundary differential equations. The particular solutions are then determined. As homogeneous solutions of the \( i \)th order (\( i = 2, 3 \)) are taken

\[
\eta_i^H = C_1 \cos \Theta_{i1} + C_2 \cos \Theta_{i2}, \quad (34)
\]

\[
\Phi_i^H = D_1 I_0(k_{1i}) \sin \Theta_{i1} + D_2 I_0(k_{2i}) \sin \Theta_{i2}, \quad (35)
\]

with \( \Theta_{i1} = 2\omega_1 t - k_{2i} z \) and \( \Theta_{i2} = 3\omega_1 t - k_{3i} z \), \( k_{2i} \) and \( k_{3i} \) coming from the dispersion relation with \( \omega = 2\omega_1 \) and \( 3\omega_1 \), respectively. The homogeneous solutions are substituted in the boundary differential equations. A relation between \( \eta_i^H \) and \( \Phi_i^H \) is deduced. The complete solution is found by substitution of the homogeneous and particular part of the solution in the nozzle conditions. The equations from which all constants, necessary for the complete solution, can be calculated are given by Busker (1983). In the case of \( \omega_1 < \omega \), the \( k_{1,2} \) values are conjugate complex and it is possible to make simplifications in \( R \) and \( \Phi \). The third-order solution is given in the appendix.

If necessary the velocity components, pressure distribution and volume in the jet could be calculated up to the breakup point. The field of interest in this paper is restricted to the
unstable growth modes of the disturbances, i.e. \( \omega_1 < \omega_\star \). Because \( R(z) \) is continuous, \( R \) cannot describe the discontinuous behaviour of the jet beyond the break-up point. Therefore only \( z \leq z_B \) is considered.

4. Application

4.1. Comparing the model with a particular experiment

Chaudhary and Maxworthy (1980a, 1980b) investigated experimentally the breakup process with a piezoelement as a disturbance source. This was driven by a sinusoidal wave at a frequency \( f^* = 100 \text{ kHz} \) with a peak voltage \( V_e^* \) as amplitude. The jet velocity \( v_0^* \) was fitted to reach the desired wavelength \( \lambda^* \). It resulted in the dimensionless wavenumber

\[
    k = \frac{2\pi r_0^*}{\lambda^*}. \tag{36}
\]

Chaudhary and Maxworthy (1980a) used a temporal instability model to determine the dimensionless minimum breakup time \( T_B^* \). In this article a relation is deduced between the theoretical disturbance amplitude \( \delta_0^* \) and the peak voltage \( V_e^* \) of their experiments. The results on the experimental break-off mode are compared with the predictions of the present model. The experimental data are transformed to desirable values of \( \omega_1 \), We and \( z_B \) for the spatial instability model with the relations

\[
    v_0^* = \lambda^* f^*, \tag{37}
\]

\[
    \text{We} = \frac{2\pi^2 \rho^* r_0^* (f^*)^2}{\sigma^* (k^*)^2} \tag{38}
\]

and

\[
    z_B^* = (\text{We} T_B^*)^{1/2}. \tag{39}
\]

The theoretical value of \( \omega_1 \) is found by a variation of \( \omega \) in the dispersion relation at the given Weber number to effect that \( \text{Re}\{k_{1,2}\} = k \). The Reynolds number in the jet is given by

\[
    \text{Re} = \frac{\rho^* v_0^* r_0^*}{\mu^*} \approx \frac{2\pi p^* (r_0^*)^2 f^*}{\mu^* k}. \tag{40}
\]

Chaudhary and Maxworthy (1980b) used a water jet and varied \( k \) between 0.3 and 1. The values of the physical quantities are given in table 1. This resulted in values of We in the range of 170 to 1910 and values of Re in the range of 640 to 2140. A series of experimental data has been compared with a series of numerical results of our model. The wavenumber \( k = 0.4312 \) is used which gives \( \text{We} = 922.6, \text{Re} = 1486 \) and \( \omega_1 = 0.4313 \). Some pictures are given by Chaudhary (1980a). In case of \( \text{We} = 922.6 \) the experimental amplitude \( V_e^* \) varied between 1 and 80 V and the corresponding minimum breakup length \( z_B \) between 860 and 305. A breakup length smaller than 586 causes a satellite breakup after the main drop, as opposed to the breakup before the

| \( n_0^* \) (m) | 3.048 \times 10^{-5} |
| \( \sigma^* \) (N/m) | 65.3 \times 10^{-3} |
| \( \rho^* \) (kg/m\(^3\)) | 1002 |
| \( \mu^* \) (Pa s) | 9.128 \times 10^{-4} |
main drop for higher values of $z_B$. The corresponding calculated values of $z_B$ lead to a range of dimensionless disturbance-amplitude values $\delta_0$ between $3 \times 10^{-3}$ and $4 \times 10^{-2}$ (table 2).

$V_\star^*$ and $\delta_0$ can be written in a logarithmic linear relation. The constants are calculated by linear regression leading to

$$\ln(V_\star^*) - 9.26 + 0.896 \ln(\delta_0).$$

A second series of experimental data with $k = 0.720$ and $We = 330.9$ is used as a comparison. The experimental values of $z_B$ were between 417 and 169. The calculated values of $\delta_0$ with $\omega_1 = 0.720$ are in the range between $3 \times 10^{-5}$ and $3 \times 10^{-3}$ (table 3).

The same log-linear relation for $V_\star^*$ and $\delta_0$ does not apply when breakup lengths are smaller than 300, see fig. 3.

### 4.2. Breakup mode of main drops and satellites

Another test criterion for the validity of the model is the break-off mode of main drops and satellites at the minimum breakup length $z_{B_{\min}}$. For small amplitudes $\delta_0$ the main drop will break off first of all at the breakup length $z_B = z_{B_{\min}}$ followed by the satellite at $z_B > z_{B_{\min}}$.
Fig. 3. Comparison between calculated values of the breakup length $z_b$ (+) with experimental data from Chaudhary and Maxworthy (1980a, 1980b) (○). Curve (1) represents values of $z_b$ in case of $We = 922.4$ and $\omega_1 = 0.4313$. Curve (2) in case of $We = 330.9$ and $\omega_1 = 0.720$ under the condition of eq. (41).

Fig. 4. The calculated jet shape just before the breakup point $z_b$ for small values of the disturbance amplitude $\delta_0$. The calculated diameters and their lengths of the main drop and satellite are given.

Fig. 5. The calculated jet shape just before the breakup point $z_b$ for large values of the disturbance amplitude $\delta_0$. The calculated diameters and the lengths of the main drop and satellite are given.
For large amplitudes $\delta_0$ the satellite will break off first of all at $z_B = z_{B\text{min}}$ followed by the main drop. The jet shape in the surroundings of the breakup length $z_B$ for small and large $\delta_0$ values are shown in figs. 4 and 5. In the case of $\omega_1 = 0.4313$, $We = 922.6$ the experimental transition value of the different modes occur for $V_e^*$ between 4 and 10 V (Chaudhary and Maxworthy, 1980a) corresponding with $\delta_0$ between $15 \times 10^{-6}$ and $4 \times 10^{-4}$ for the fixed log-linear relation between $V_e^*$ and $\delta_0$. The analytical model predicts that the transition value for $\delta_0$ lies between $5 \times 10^{-3}$ and $6 \times 10^{-3}$. This discrepancy is probably due to the neglect of viscous effects in the jet ($Re = 1486$). The reason for this is the slow specific velocity of the breakup process $(\sigma^*/\rho^* r_0^*)^{1/2}$, in which the viscosity effects can not be neglected. The model of Chaudhary and Maxworthy (1980a) seems to predict the break-off mode.

4.3. The forming of satellites

Sterling and Sleicher (1975) demonstrated the unstable breakup for viscous liquid jets with aerodynamic effects. Their model and the model described in this article do not predict the forming of satellites for $\omega_1 \geq 0.7$. Lafrance (1974), Rutland and Jameson (1970) did not succeed in finding satellites either. Fig. 6 shows the divergence between model and experiments.

Fig. 6. Calculated jet shape at $We = 300$, $\omega = 0.7$, $\delta_0 = 1 \times 10^{-4}$ shows no forming of satellites.

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Fig. 7. The forming of two satellites under the conditions $We = 300$ and $\delta_0 = 1 \times 10^{-2}$ in case the dimensionless frequency $\omega_1 = 0.3$. The calculated wavelength is given.
4.4. Effect of the disturbance amplitude on the dimensions of satellites and the main drop

The volume of satellites and main drops is dependent on the disturbance amplitude $\delta_0$. The ratio of the maximum diameters of satellites and main drops changes the same way as the ratio of their lengths changes with $\delta_0$. These phenomena can be seen in figs. 4 and 5. The ratio of the diameters of the satellite to the main drop for $\delta_0 = 1 \times 10^{-2}$ is 0.579 while the ratio of the lengths is 1.05. In case of $\delta_0 = 1 \times 10^{-6}$ the ratio of the diameters becomes 0.479 and the ratio of the lengths 0.728. This means that satellites are larger for greater values of $\delta_0$. About this no experimental data are available. Because of the limited use of this model, this point will not be pursued further.

4.5. Breakup length as function of the frequency

The dependence of the minimum breakup length $z_B$ on the frequency for several Weber and amplitude numbers is shown in fig. 8. The plot shows that the minimum breakup length over the whole range $0 < \omega_1 < \omega_S$ occurs somewhere between $\omega_1 = 0.7$ and $\omega_1 = 0.8$ for small $\delta_0$, and moves towards $\omega_S$ for large $\delta_0$. It can be expected that an undisturbed jet picking up a disturbance from its surroundings will not break up at the most unstable mode (just below $\omega_1 = 0.7$) but that it breaks up at frequencies in an interval $0.65 < \omega_1 < 0.8$.

5. Discussion and conclusions

The theoretical model of spatial instability provides a good insight into the breakup process in a jet and may be extended. In some respects the model is in accordance with the results of
experiments published in literature, in other respects it show discrepancies therewith. The model describes the physical reality better than models based on temporal instability. Our model and Yuen's (1968) are analogous and have as common features the forming of two satellites for $\omega_1 < 1/3$ and the absence of satellites for $\omega_1 > 7/10$.

The model is more extensive than the one-dimensional Cosserat theory because we made use of two-dimensional fluid mechanical equations taking account of the flow in radial direction. However, by comparing the theoretical predictions with available experiments data, it follows that the discrepancies exist due to the neglect of the viscosity or higher-order terms of the solution. Further research could take into account the effect of the nozzle shape and the disturbances on the breakup mode of the jet. In addition, the effects of gravity, aerodynamical, mass and heat transfer phenomena could be investigated.

Appendix

The third-order solution for the jet surface is

$$R = 1 + 2\delta_0 P_{1,1} \text{Im}\{\cos \Theta_1\}$$

$$+ \delta_0^2 \text{Re}\{P_{2,1} \cos \Theta_2 + P_{2,2} \cos \Theta_3 + 2P_{2,3} \cos 2\Theta_1 + P_{2,5} \cos[2(\omega_1 t - \text{Re}\{k\} z)]$$

$$+ P_{2,6} \left[\cosh(2\text{Im}\{k\} z) - 1\right]\}$$

$$+ \delta_0^3 \text{Im}\left\{2 \sum_{i=1}^2 P_{3,i} \cos \Theta_3 + 2 \sum_{i=1}^2 \left[P_{3,4i} \cos(\Theta_2 + \Theta_1) + P_{3,4i+1} \cos(\Theta_2 - \Theta_1)\right]\right\}$$

$$+ 2P_{3,12} \cos(3\Theta_1) + 2P_{3,14} \cos(2\Theta_1 + \Theta_2) + 2P_{3,15} \cos(2\Theta_1 - \Theta_2)$$

$$+ (2P_{3,18} + P_{3,20}) \cos \Theta_4 + P_{3,21} \cos \Theta_2$$

with $\Theta_1 = \omega_1 t - k_1 z$ and $\Theta_2 = \omega_2 t - k_2 z$. The subscripts $i$ and $j$ of the constants $P_{i,j}$ indicate the $i$th order and the number of the constant respectively.

To obtain numerical results for the analytical solution of the jet surface it is necessary to calculate the $k$ values for the first, second and third order using $\omega_1$ and the Weber number as parameters. The calculation of the various constants $P_{i,j}$ is the next step. After that the jet surface $R$ can be determined for different amplitudes $\delta_0$ and time $t$ in a given distance interval $\Delta z$. From these results computer plots can be made. A description of the software is given by Busker (1983).

References


