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Finite thermal diffusivity at onset of convection in autocatalytic systems: Discontinuous fluid density

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A linear convective stability analysis for propagating autocatalytic reaction fronts includes density differences due to both thermal and chemical gradients. Critical parameters for the onset of convection are calculated for an unbounded geometry, a vertical slab, and a vertical cylinder. Thermal effects are important at unstable wavelengths well above the critical wavelength for the onset of convection. © 1995 American Institute of Physics.

Experiments observe convection for ascending iodate-arsenous acid fronts in vertical tubes whose diameters exceed a critical diameter of about 1 mm. Linear hydrodynamic stability theory predicts a transition to buoyancy-driven convection with a critical diameter that agrees with these experiments, although large experimental and theoretical uncertainties preclude a precise comparison. A 24% theoretical uncertainty is due to the difference between the critical diameters calculated for infinite and zero thermal diffusivity. To reduce this uncertainty, calculations are needed at finite thermal diffusivity.

The equation of state \( p(T) = p_1 \gamma[1 - \alpha(T - T_f)] \) is a first-order Taylor expansion of the temperature-dependent mass density about \( T_f \), the final adiabatic temperature of the reacted fluid. Here, \( \gamma = 1 \) in the unreacted fluid and \( \gamma = 1 - \delta \), in the reacted fluid to account for chemically driven density changes, and \( \alpha \) is the isobaric thermal expansion coefficient.

Neglecting small density changes \( \delta \approx \Delta T \approx 1 \times 10^{-4} \) except where they modify gravity to first order allows us to write the dimensionless thermohydrodynamic equations of motion as

\[
\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla) \mathbf{v} = -\nabla \mathbf{P} + \frac{\mathbf{g}}{\rho_0} \mathbf{g}, \quad (1a)
\]

\[
\mathbf{v} \cdot \mathbf{n} = 0, \quad (1b)
\]

and

\[
\frac{\partial T}{\partial t} + \mathbf{v} \cdot \nabla T = \frac{\mathbf{g}}{\rho_0} \mathbf{g} \nabla^2 T. \quad (1c)
\]

These are measured in units of length \( a \) to be specified below, time \( a^2/v \), pressure \( \rho_1 v^2/a^2 \), and temperature \( a \Delta T / \alpha \), where \( \mathbf{v} \) is the kinematic viscosity. The convenient reduced pressure \( P \) is related to the conventional pressure by \( P = P_{\text{conv}} + (\gamma \mathbf{g} a^2/v^2) z \). The equations also involve a dimensionless thermal diffusivity \( \mathcal{D}_T = D_T/v \), a dimensionless molecular diffusivity \( \mathcal{D}_C = D_C/v \), and a dimensionless driving parameter \( \mathcal{F} = \delta a^2 g/v D_C \).

The location of the thin chemical front is given by \( z = h(x, y, t) \), a function of the horizontal coordinates and the time, obeying the eikonal relation

\[
\hat{n} \cdot \frac{\partial z}{\partial t} = v_0 + \mathcal{D}_C \kappa + \hat{n} \cdot \mathbf{v} |_{h}. \quad (1d)
\]

Here, \( \hat{n} \) is the normal vector to the front pointing into the unreacted fluid, \( v_0 = (a/\mathbf{v}) c_0 \) is the dimensionless flat front speed, \( \mathbf{v} |_{h} \) is the fluid velocity evaluated at \( z = h(x, y, t) \), and \( \kappa \) is the curvature of the front measured as positive when the center of curvature is in the unreacted fluid.

The boundary conditions require that the velocity and the normal derivative of the temperature vanish at the sidewalls, ensuring no-slip, insulating sidewalls. The jump conditions at the front express the conservation of energy, mass and momentum:

\[
\begin{align*}
[\hat{n} \cdot \mathbf{v}]_{h} &= 0, \quad (1e) \\
[\hat{n} \cdot \mathbf{v}]_{h} &= 0, \quad (1f) \\
[e_{ijk} n_i n_j T_{u}^{i}]_{h} &= 0, \quad (1g) \\
[P]_{h} &= -\mathcal{D}_C h - [n_i n_j T_{u}^{i}]_{h}, \quad (1h) \\
[\hat{n} \cdot \mathbf{v} T]_{h} &= \frac{\Delta T}{\partial h} \left( \frac{\partial h}{\partial t} - \hat{n} \cdot \mathbf{v} |_{h} \right), \quad (1i)
\end{align*}
\]

and

\[
[T]_{h} = 0. \quad (1j)
\]

Here, \( [q]_{h} \) denotes the jump in a given function \( q - q(x, t) \) in the following way. If \( q(x, t) = q_{r}(x, t) \) in the reacted fluid and \( q(x, t) = q_{u}(x, t) \) in the unreacted fluid, then \( [q]_{h} = q_{r} |_{h} - q_{u} |_{h} \) is the difference between these functions evaluated at the reaction front \( z = h \). Furthermore, \( T_{ij} = -\partial v_{i}/\partial x_{j} \) is the (dimensionless) viscous stress tensor, \( e_{ijk} \) is the totally antisymmetric tensor, \( n_{i} \) are the Cartesian components of \( \mathbf{n} \), and \( \Delta T = T_0 - T_f \) involves the initial fluid temperature \( T_0 \). Equations (1) determine the convective motion and stability of a thin autocatalytic reaction front.

Steady upward propagation of a convectionless flat front is conveniently described in a frame of reference that is stationary with respect to the front. With the front located at

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and

\[ T_u^{(0)} - T_0 \pm \Delta T \exp \left( \frac{-z_{0\gamma}}{D_T} \right) \]

in the reacted and unreacted fluids. This solution lacks both fluid motion (in the laboratory frame) and horizontal thermal gradients. Hence, it is applicable to both unbounded geometries and geometries with vertical insulating sidewalls.

To study the stability of the convectionless flat front, we perturb it slightly according to \( \mathbf{v} = \mathbf{v}^{(0)} + \mathbf{v}^{(1)} \), with similar forms for the other dependent variables. In evaluating Eqs. (1) for small \( h \), it is convenient to employ Taylor expansions,

\[ q_h = q_0 + h \partial_z q_0 + \frac{1}{2} h^2 \partial_z^2 q_0 + \cdots \]

and

\[ [q]_h = [q_0] + h [\partial_z q_0] + \frac{1}{2} h^2 [\partial_z^2 q_0] + \cdots , \]

where \( \partial_z = \partial / \partial z \). Note that even though \( q(x,t) = q_0(x,t) \) only for \( z > h(x,y,t) \) and \( q(x,t) = q_1(x,t) \) only for \( z < h(x,y,t) \), the functions \( q_0(x,t) \) and \( q_1(x,t) \) are defined over all space, so that the difference \( [q]_0 = q_0 - q_1 \) between these functions at \( z = 0 \) is a well-defined quantity. Using these in Eqs. (1) and linearizing in the perturbations yields

\[
\begin{align*}
(\partial_t - v_0 \partial_z - \nabla^2) \mathbf{v}^{(1)} &= \mathbf{P}_T \mathbf{v}^{(1)} - \nabla P^{(1)}, \\
(\partial_t - v_0 \partial_z - \nabla^2) T^{(1)} + (\partial_z T^{(0)}) w^{(1)} &= 0, \\
\nabla \cdot \mathbf{v}^{(1)} &= 0, \\
(\partial_t - \mathbf{P}_T \nabla^2) h^{(1)} &= w^{(1)} |_{0}, \\
[w^{(1)}]_0 &= 0, \\
[z \times \mathbf{v}^{(1)}]_0 &= 0, \\
[\epsilon_{3h} k^3_{3h}]_0 &= 0, \\
[P^{(1)}]_0 &= - \mathbf{P}_T h^{(1)} + 2[\partial_z w^{(1)}]_0, \\
[T^{(1)}]_0 &= - \frac{\Delta T v_0}{D_T} h^{(1)},
\end{align*}
\]

and

\[ [\partial_z T^{(1)}]_0 = \frac{\Delta T}{D_T} \left( \frac{v_0^2}{D_T} + \mathbf{P}_T \nabla^2 \right) h^{(1)}. \]

Here, we have appropriately set \( \mathbf{v}^{(1)} = - \nabla h^{(1)} \) (a horizontal vector) and \( \mathbf{P}^{(1)} = \nabla^2 h^{(1)} \), and have written the vertical velocity as \( w^{(1)} = - \mathbf{v}^{(1)} \).

The unbounded system is the simplest possible geometry in which to study convection in autocatalytic systems. Here we analyze the time dependence of infinitesimal perturbations about a convectionless flat unbounded front. For simplicity, we analyze two-dimensional perturbations \( \mathbf{v}^{(1)}(x,z,t) = h^{(1)}(x,z,t) \hat{x} + w^{(1)}(x,z,t) \hat{z} \) by assigning the exponential dependence \( e^{i(\mathbf{k} \cdot \mathbf{x}) + \sigma t} \) to all perturbed quantities. Substituting these forms into Eqs. (4) yields a system of ordinary differential equations that is similar to Eqs. (9) in Ref. 5, except for an important contribution from the discontinuous density. Exponential solutions are readily obtained after the manner of Ref. 5.

Figure 1 shows the resulting growth rate \( \sigma \) of infinitesimal perturbations about a flat front as a function of the wave number \( q \) for the typical parameters \( D_C = 2 \times 10^{-3} \) cm²/s, \( \mathbf{v} = 0.87 \times 10^{-4} \), \( \Delta T = 0.4 \) °C, \( \nu = 9.2 \times 10^{-3} \) cm²/s, and \( c_0 = 2.95 \times 10^{-3} \) cm/s for various thermal diffusivities \( D_T \), measured in units of cm²/s, with the dashed trace corresponding to the thermal diffusivity of water.
$d > d_T$, the thermally induced density changes occur over a region small compared with the convective region, which extends well into the unreacted fluid at temperature $T_0$. The crossover between zero and infinite diffusivity should therefore occur when $d \approx d_T$, that is, when $D_T^{-1} \approx 2\pi n_0 q$. Since the values $11 < q_m < 17$ of $q_m$ are much smaller than the values $150 < q_c < 200$ of $q_c$ over the whole range of $D_T$ (Fig. 2), the crossover value $D_T^{-1} \approx 4 \times 10^{-3}$ cm/s for $q_m$ must exceed the crossover value $D_T^{-1} \approx 3 \times 10^{-4}$ cm/s for $q_c$. Indeed, these crossover values are consistent with the crossover behavior in Fig. 2.

Intriguing minima in $q_c$ and $q_m$ occur in Fig. 2 near the crossover values of $D_T$, where the convective and thermal length scales are comparable. At the minimum in $q_c$, the flat front is most stable to convective perturbations, and is even more stable than for infinite thermal diffusivity. This is due to thermal stabilization of the perturbed front: Because heat flow is normal to the front surface [Eq. (11)], the unreacted fluid above the valleys in the front is heated faster than the fluid above the peaks. Thermal expansion consequently renders the fluid above the valleys more buoyant than otherwise, thus tending to raise the valleys and lower the peaks. This stabilizing mechanism is absent at infinite thermal diffusivity where the fluid temperature is effectively uniform throughout. It is also absent at zero thermal diffusivity, where all thermal gradients are confined to the front surface itself. Thermal stabilization is only relevant when the convective and thermal length scales are comparable, where it produces the minima in $q_c$ and $q_m$.

The two-dimensional motion of a fluid confined between two parallel vertical planes located at $x = \pm a/2$ and the axisymmetric motion of a fluid confined to the interior of a long vertical cylinder of radius $a$ are studied by generalizing the expansion procedure in Ref. 4. In each case, a flat ascending front is linearly unstable to convection above a critical length scale $a_r$ for the onset of convection (Fig. 3). The limits of zero and infinite thermal diffusivity agree with our previous calculations. The maximum in $a_r$ implies maximum stability, and corresponds to the minimum in $q_c$ discussed earlier. Results for $D_T = 1.45 \times 10^{-3}$ cm/s (vertical arrow) can clearly be approximated by the infinite diffusivity limit.

In conclusion, chemical gradients overwhelm thermal gradients in determining the length scale for the onset of convection in autocatalytic reaction fronts. However, at larger convective length scales approaching the thermal length scale, thermal gradients are important. An additional mechanism for stabilizing ascending reaction fronts, relevant only at finite thermal diffusivity, arises from the enhanced (depressed) heating rates above valleys (peaks) in the front, which renders the fluid above the valleys hotter and more buoyant than otherwise. This mechanism may play a crucial role in systems for which chemical and thermal gradients compete rather than cooperate with each other to destabilize the system.

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