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Front propagation in cellular flows for fast reaction and small diffusivity

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We investigate the influence of fluid flows on the propagation of chemical fronts arising in Fisher-Kolmogorov-Petrovsky-Piskunov (FKPP) type models. We develop an asymptotic theory for the front speed in a cellular flow in the limit of small molecular diffusivity and fast reaction, i.e., large Péclet (Pe) and Damköhler (Da) numbers. The front speed is expressed in terms of a periodic path—an instanton—that minimizes a certain functional. This leads to a compact expression for the speed in the limit of slow reactions and small diffusivity [16,17]. Here we consider the opposite limit of fast reaction (e.g., Ref. [18]) as asymptotic regimes. Results of this type have been derived in [6–8], with much effort devoted to steady spatially periodic flows such as (2), this is given by

\[ \psi(x, y) = -\sin x \sin y. \]  

We take the domain to be the strip \( D = (-\infty, \infty) \times [0, \pi] \) along which an infinite array of identical cells are arranged, each of which is composed of two half cells of opposite circulation with hyperbolic stagnation points at each corner (streamlines are shown for a single cell in Fig. 2). As the initial condition we take \( \theta(x, y, 0) = \theta(-x) \), where \( \theta \) is the Heaviside step function. The boundary conditions are \( \theta(\infty, y, t) = 0 \) and \( \theta(-\infty, y, t) = 1 \), so that the front advances rightwards, and no flux, \( \partial_x \theta = 0 \) at \( y = 0, \pi \). The front characteristics change drastically with Da. When Da is small, the front’s leading edge is confined near the cell boundaries (see left column in Fig. 1). As Da increases, the front sharpens and the imprint of the flow, the boundary-layer structure in particular, is less prominent (see middle and right columns in Fig. 1).

The long-time speed of propagation of the front is determined by the behavior of the solution near the front’s leading edge, where \( 0 < \theta \ll 1 \) and \( r(\theta) \approx r'(0) = \theta \). For steady periodic flows such as (2), this is given by

\[ e = \inf_{q > 0} \frac{f(q) + Da}{q}, \]  

where \( f(q) \) is the largest eigenvalue of

\[ \mathcal{L}v = f(q)v, \]  

with the operator \( \mathcal{L} \) defined by

\[ \mathcal{L} = Pe^{-1} \Delta - u \cdot \nabla - 2 Pe^{-1} q \partial_x + (u_1 q + Pe^{-1} q^2), \]  

and periodic and no-flux boundary conditions in \( x \) and \( y \), respectively (see Refs. [15,20 (Chap. 7),21]). This result is intimately connected with the long-time large-deviation rate
In the absence of a background flow, \( f(q) = q^2/Pe \), recovering the classical formula for the bare speed \( c_0 = 2\sqrt{Da/Pe} = 2\sqrt{\kappa/\tau} \) (see Refs. [13,14,24,25] and references therein). For general \( u \neq 0 \), the eigenvalue problem (4) cannot be solved explicitly.

Small diffusivity, fast reaction. Our purpose is to use asymptotic analysis to obtain the speed of the front in the large-Péclet limit with

\[
\gamma = Da/Pe = O(1),
\]

corresponding to the geometric optics regime defined by \( \kappa, \tau \to 0 \) with \( \kappa/\tau = O(1) \) [26]. This can be achieved by analyzing the large-Péclet limit of the eigenvalue problem (4), or by considering the large-\( t \) limit of the geometric optics theory treatment of [20 (Chap. 6)] (see also Refs. [18,27]). It is, however, convenient to start with the basic result exploited by Refs. [15,20], namely, that the front speed is controlled by the point at which the solution to the linearization of Eq. (1) neither grows nor decays.

We seek a solution to this equation in the WKB form

\[
\theta(x,y,t) \sim \exp[-Pe[I(x,y,t) - \gamma t]],
\]

for \( Pe \gg 1 \), (5) where \( I(x,y,t) \) can be recognized as the small-noise large-deviation rate function [28]. At leading order, \( \delta I + H(\nabla I, x, y) = 0 \), where \( H = \|\nabla I\|^2 + u(x,y) \cdot \nabla I \) is the Hamiltonian and \( \|\cdot\| \) the usual norm. Its solution is well known from Hamilton-Jacobi theory (e.g., Refs. [28,29]) and is given by

\[
I(x,y,t) = \frac{1}{4} \inf_{\phi(t)} \int_0^t \|\phi(s) - u(\phi(s))\|^2 ds,
\]

subject to \( \phi(t) = (x,y) \) and \( \phi(0) = (0,.) \). Thus, the behavior of (5) is controlled by a single path \( \phi^*(t) \) that minimizes this integral. This optimal path is often called instanton [30]. Expression (5) then indicates that for \( t \gg 1 \), the front speed satisfies

\[
\gamma = \lim_{t \to \infty} \frac{I(x(t),y,t)}{t} \equiv \mathcal{G}(c),
\]

where the limit \( t \to \infty \) eliminates the dependence on \( y \). This leads to

\[
c = \mathcal{G}^{-1}(\gamma). \tag{8}
\]

The front speed is therefore obtained by calculating \( \mathcal{G}(c) \). This calculation is significantly simplified by observing that the limit in (7) is determined in terms of periodic trajectories, an observation justified rigorously in Ref. [31]. We take solutions \( \phi(s) \) to be periodic in the sense that \( \phi(t) = \phi(t + 2\pi/s) \), where the period is \( s = 2\pi/c \). Letting \( \sigma = cs \), we obtain the simplified expression

\[
\mathcal{G}(c) = \frac{1}{8\pi} \inf_{\phi(t)} \int_0^{2\pi} \|c\phi'(\sigma) - u(\phi(\sigma))\|^2 d\sigma, \tag{9}
\]

subject to \( \phi(2\pi) = \phi(0) + (2\pi,0) \). Expressions (8) and (9) are the main result of this Rapid Communication. They provide a direct way of computing the instantons and thus the front speed. Note that \( \mathcal{G}(c) \) may be interpreted as the Legendre transform of \( H\overline{H} \), the effective Hamiltonian of the homogenized Hamilton-Jacobi equation \( \delta I + H(\nabla I) = 0 \) [18,32]. In order to derive (8), we have formally assumed that \( \gamma = O(1) \). However, the asymptotic results apply over a broad range of values of \( \gamma \), specifically \( \gamma \gg (Pe \log Pe)^{-1} \), as we show below.

The solution to the minimization problem (9) is easy to obtain numerically. We first discretize (9) and use MATLAB to minimize the action. We iterate over \( c \), starting with large
values for which the straight line $\phi^*(\sigma) = (\sigma, \pi/2)$ is a good initial guess. Characteristic examples of instants obtained for different values of $c$ are shown in Fig. 2 (the two cases with $c = 0.5$ and $c = 1$ correspond to the first two columns in Fig. 1). For large values of $c$, the instanton is close to a straight line. For small values of $c$, it follows closely a streamline near the cell boundaries. Figure 3 shows the behavior of $c$ as a function of $c$ deduced from (8).

Asymptotic limits. We now obtain closed-form expressions for the speed $c$ in two asymptotic limits. The first one corresponds to $c \ll 1$ and hence $c \ll 1$. Numerical results (see Fig. 2 for $c = 0.5$ and $c = 0.5$) suggest that the instanton departs from the streamline only for $c \approx \pi/2$ when it crosses the separatrices between adjacent cells. It is also clear that (9) is minimized when $\phi^*(\sigma) = (x(\sigma), y(\sigma))$ satisfies $c y' \approx -\cos x \sin y$ (so that the instanton and flow speeds differ only in the $x$ direction).

Exploiting symmetry to consider $0 \leq \sigma \leq \pi/2$ only, with $x(0) = 0$, $y(0) = x(\pi/2) = \pi/2$, and $y'(\pi/2) = 0$, we divide the instanton into two regions (see Fig. 2). In region 1, where $x \ll 1$, the integrand in (9) is approximately $(c x' - x \cos y)^2$, leading to the Euler-Lagrange equation $c^2 x'' = x$ (since $c y' \approx -\sin y$). In region 2, $y \ll 1$, $c x' \sin x \cos y \approx -\sin x$, and $c y' = -\cos x \sin y \approx -c \cos x$. Matching the solutions for $x,y \ll 1$ (the cell corner) gives the approximation

$$\phi^*(\sigma) \approx \begin{cases} (C_1(\sigma), C_2(\sigma)) & \text{for } \sigma \ll \pi/2, \\ (C_2(\pi/2 - \sigma), C_3(\pi/2 - \sigma)) & \text{for } \sigma \gg c. \end{cases}$$

Expression (10) is in very good agreement with our numerical solution (see Fig. 2 for $c = 0.5$).

Using (10) gives the integrand in (9) as $(c x' - x \cos y)^2 \approx 16 \exp(-\pi/c) \cosh^{-1}(\sigma/c)$, leading to

$$\mathcal{J}(c) \sim 4 \times (2/\pi) e^{-\pi/c}, \quad \text{where } c \ll 1,$$

and the factor 4 appears because, for $\sigma \in [0, 2\pi]$, there are four regions that are similar to region 1. Inverting (11) and using (8) finally gives

$$c \sim \frac{\pi}{W_p(8\gamma^{-1})}, \quad \text{for } \gamma \ll 1,$$

where $W_p$ denotes the principal real branch of the Lambert $W$ function [33]. This approximation holds for $(\log Pe^{-1}) \ll Da \ll Pe$: the lower bound follows from requiring that the argument of the exponential in (5) be large, $Pe I \sim Da \tau \gg 1$, where the period is roughly estimated as $\tau = 2\pi/c \approx \log Pe/2$ (see (13) below and Ref. [23] for a complete argument). Figure 3 shows that this approximation is excellent within its range of validity. Since $\gamma \ll 1$, it is consistent to approximate $W_p(8\gamma^{-1})$ to reduce (12) to

$$c \approx \frac{\pi}{\log Pe}.$$

A qualitatively similar expression was obtained in Refs. [34,35] using a heuristic approach based on the so-called $G$ equation. Note that the logarithmic dependence of the speed on $Pe$ and its slow growth with $Da$ [captured by (12) but not (13)] is associated with the holdup of the instanton near the hyperbolic stagnation points at the cell corners.

The second limit leading to closed-form results corresponds to $c \gg 1$, hence $c \gg 1$. In this case, we seek an instanton as a power series $\phi^*(\sigma) = (\sigma, y_0) + c^{-1}(x_1(\sigma), y_1(\sigma)) + \cdots$, where $x_1, y_1$ are functions of period $2\pi$ that satisfy $x_1(0) = y_1(0) = 0$. Substituting into (9), we find that at $O(c^{-1})$,

$$\mathcal{J}(c) \approx \frac{c^4}{4} + \frac{1}{8\sigma} \inf_{x_1, y_1} \int_0^{2\pi} \left(x_1^2 + y_1^2 + 4 y_1 \sin \sigma \sin y_1\right) d\sigma,$$

(14)
after some manipulations. Minimizing this integral leads to the instanton

\[ \phi^*(\sigma) = \left( \sigma, \frac{\pi}{2} \right) + c^{-1}(0, -2 \sin \sigma) + \cdots, \quad \text{for } c \gg 1, \quad (15) \]

in excellent agreement with our numerical solution (see Fig. 2 for \( c = 5 \)). Combining (14) and (15) yields \( \phi(c) = c^2/4 - 3/8 + O(c^{-2}) \). We now use (8) to find that

\[ c \sim 2\sqrt{\gamma} \left( 1 + \frac{3}{16\gamma} + \cdots \right), \quad \text{for } \gamma \gg 1, \quad (16) \]

which corresponds to \( Da \gg Pe \). The leading-order term in (16) is the bare speed \( c_0 \). As Fig. 3 shows, the second term in the expansion is necessary for a good agreement between asymptotic and full results.

**Comparison with numerical results.** We now compare our predictions for \( c \) derived from (8)–(9) with values obtained from (i) a numerical evaluation of the principal eigenvalue in (4), and (ii) direct numerical simulations of (1) with \( r(\theta) = \theta/(1 - \theta) \). For (i) we use a standard second-order discretization to approximate (4) and choose the spatial resolution \( \Delta \) to satisfy \( \pi/\Delta = 750 \). The resulting matrix eigenvalue problem is solved for a range of values of \( q \) using MATLAB. For (ii) we discretize (1) using a fractional-step method with a Godunov splitting which alternates between advection (using a first-order upwind method with a minmod limiter—see Ref. [36] for details), diffusion (using an alternating direction implicit method), and reaction (solved exactly). We choose the same spatial resolution \( \Delta \) as for (4). The computational domain is made finite using artificial boundaries at \( x = \pm N\pi \), with \( N = 5 \), so that boundary effects are negligible. The front is tracked for long times by modifying the computational domain: When the solution at \( x = (N - 1)\pi \) exceeds \( \delta = 10^{-6} \), we eliminate the nodes with \( -N\pi \leq x \leq -(N + 1)\pi \) and add new nodes with \( N\pi \leq x \leq (N + 1)\pi \) where we set \( \theta = 0 \). We calculate the front speed using a linear fit of the right endpoint of the front, \( x^\ast(t) = \max\{x : \theta(x,t) = \epsilon\} \) where \( \epsilon = 10^{-3} \). Results are insensitive to the exact values of \( \epsilon \) and \( \delta \).

The three sets of numerical results are shown in Fig. 4. The speeds derived from the eigenvalue equation (4) are in excellent agreement with the corresponding values obtained from the full numerical simulations of Eq. (1). With increasing values of \( Pe \), the asymptotic expression (9) becomes increasingly accurate, with excellent agreement for \( Pe = 250,500 \) and satisfactory agreement for the moderate values \( Pe = 50,125 \). As expected, (8)–(9) is valid for a broad range of values of \( Da \), restricted only by \( Da \gg (\log Pe)^{-1} \). Note that the use of both the eigenvalue equation and the full numerical simulations is restricted: as \( Pe \) increases, the solutions to (1) and (4) become progressively localized, with \( O(1/\sqrt{Da}Pe) \) length scales that are challenging to resolve when \( Da,Pe \gg 1 \).

**Conclusion.** We have derived a compact expression for the front speed based on the minimization of the large-deviation action (9) over periodic instantons. This leads to the efficient computation of the speed for a large range of values of \( Da \). For the particular case of cellular flows, this expression provides the closed-form results (12) and (16) valid for \( (\log Pe)^{-1} \ll Da \ll Pe \) and \( Da \gg Pe \). In the first regime, the passage of the front near the stagnation points at the cell corners is shown to control the front speed; as a result this is almost insensitive to the reaction rate and depends logarithmically on the Péclet number. For \( Da = O(\log Pe)^{-1} \) and smaller, the front speed is not controlled by a single minimizing trajectory, and asymptotic solutions to the eigenvalue problem (4) must be sought by other means; this will be the subject of future work [23].

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