Theory of light-induced drift of a one-component gas in a capillary

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(submitted 15 February 1995)


The surface and collisional mechanisms of light-induced drift of a one-component gas in a capillary are investigated theoretically. These mechanisms are based on various interactions of the excited and unexcited particles with the surface of the capillary and with one another. The kinetic equations for the velocity distribution functions of the excited and unexcited particles are solved numerically by the discrete-ordinate method. The velocity profiles and the gas fluxes, averaged over the transverse cross section of the capillary, are calculated as a function of the Knudsen number and the ratio of the radiative decay rate to the intermolecular collision rate. © 1995 American Institute of Physics.

1. INTRODUCTION

A theory of light-induced free-molecular surface drift of a one-component gas in a capillary is presented in Refs. 1 and 2. It is shown that the Bennett peak and dip in the velocity distributions of these particles are deformed differently for various interactions of the excited and unexcited particles with the surface of the capillary, and they do not cancel one another. The asymmetry of the total distribution function relative to the zero value of the projection of the molecular velocity on the direction of propagation of the light wave determines the motion of the absorbing gas as a whole. We note that the possibility, in principle, of surface light-induced drift was first predicted in Ref. 3. In Ref. 4 calculations of this phenomenon were carried out in the hydrodynamic approximation. The main deficiency of the results obtained in Ref. 4 are that the light-induced slip boundary condition for the Navier–Stokes equation was derived under the assumption that the particle-velocity distribution function remains unchanged across the Knudsen layer.

In Refs. 5–8, collisional light-induced drift in a channel formed by two infinite parallel plates was calculated. This phenomenon owes its existence to the difference in intermolecular collision cross sections of the excited and unexcited particles. In Refs. 6–8, surface light-induced drift was also studied taking into account the spatial nonuniformity of the distribution function. In Ref. 6 this was done for both small and large values of the Knudsen number (Kn), in Ref. 7 for arbitrary values of the Knudsen number and in Ref. 8 for small values of Kn (Kn<0.2).

A kinetic theory of the surface and collisional mechanisms of light-induced drift of a one-component gas in a cylindrical capillary was developed in Ref. 9 for arbitrary Kn. However, the results obtained there are limited by the assumption that the radiative decay rate \( \Gamma_n \) of an excited level is small compared to the intermolecular collision rate \( \gamma_n \), i.e., \( \Gamma_n \ll \gamma_n \). Moreover, the Bubnov–Galerkin method used in Ref. 9 to solve the integral–moment equations for the partial velocities and viscous stresses did not make it possible to investigate the light-induced drift velocity profile and its evolution as the flow regime changed from hydrodynamic to free-molecular.

Our objective in the present work is to develop a kinetic theory of the surface and collisional mechanisms of light-induced drift of a one-component gas in a capillary for arbitrary values of Kn and the frequency parameter \( \Gamma_n^\omega \), and to calculate the velocity profile and the mean gas flux over the cross section of the capillary. This multiparameter, mathematically complicated problem can probably only be solved numerically. The optimal method appears to be the method of discrete ordinates, which is based on replacing the continuous six-dimensional phase space by a discrete space. Numerical results obtained with guaranteed accuracy are also of particular interest in that they can be used to test various approximate theories of light-induced drift.

2. ASSUMPTIONS AND BASIC EQUATIONS

We study the steady motion of a one-component gas in a circular capillary under the action of resonant optical radiation propagating along the \( z \) axis of the capillary. Let the length \( L \) of the capillary be much greater than its radius \( R_c \), so that distortions of the velocity field at the ends of the capillary can be neglected. The radiation is absorbed by the gas particles in electronic or vibrational–rotational transitions from the ground state \( n \) to an excited state \( m \). The frequency \( \omega \) of the monochromatic radiation is offset from the center \( \omega_{\text{abs}} \) of the absorption line by \( \Omega = (\omega - \omega_{\text{abs}}) \ll \omega_{\text{abs}} \). As a result of the Doppler effect, only particles with velocities \( v \) such that \( kv = \Omega \), where \( k \) is the wave vector, interact most efficiently with the radiation. Particles that have absorbed radiation have a different collision cross section. Thus, the absorbing gas can be regarded as a binary gaseous mixture in which the particles have the same masses but different interaction cross sections. Particle exchange between components of the mixture is possible as a result of radiative decay of an excited level, as well as collisional and stimulated transitions.

The velocity distributions \( f_n \) and \( f_{\text{exc}} \) of the excited and unexcited particles, respectively, have a peak and a Bennett dip, respectively, near the resonance values \( v_z = \Omega/k_z \), where \( v_z \) is the projection of the velocity vector \( v \) on the \( z \) axis. For \( \Omega \neq 0 \), these distributions are asymmetric with respect to \( v_z = 0 \). Consequently, oppositely directed macro-
scopic fluxes $J_m$ and $J_n$ of excited and unexcited particles, respectively, exist along the capillary. If the probabilities that excited and unexcited particles pass through the capillaries are different because the interactions of the two types of particles with the surface and the transport cross sections are different, then there exists a resulting flux $J = J_m + J_n$ - the light-induced drift.

In the two-level approximation, the distribution functions $f_m$ and $f_n$ satisfy the following system of kinetic equations:

$$\frac{\partial}{\partial t} f_m = \frac{1}{2} \kappa(v) (f_n - f_m) - \Gamma_m f_m + S_m,$$

$$\frac{\partial}{\partial t} f_n = \frac{1}{2} \kappa(v) (f_m - f_n) + \Gamma_n f_n + S_n,$$

where

$$\kappa(v) = \frac{4|S_m|^2}{\Gamma_m^2 + (\Delta v)^2}, \quad S_m = \frac{E_m}{2h},$$

$\Gamma_m$ is the radiative decay constant, $\Gamma$ is the homogeneous half-width of the absorption line, $S_m$ and $S_n$ are the Boltzmann collision integrals, $E_m$ is the amplitude of the electric field, $d_m = 0$ is the transition dipole moment for the transition $m \rightarrow n$, $\hbar$ is Planck's constant, and $\kappa(v)$ is the saturation parameter characterizing the probability of stimulated transitions and is proportional to the radiative intensity $I$.

Let the gas and the capillary have the same coordinate-independent temperature $T$. We also assume that collisions of molecules with the capillary surface are elastic. Then Maxwell's specular diffusion model can be chosen to give the boundary conditions for Eqs. (1). According to this model, a fraction $c_j$ of the particles in state $j$ after a collision with the surface spreads out diffusively with Maxwell's velocity distribution $f_j$, while the fraction $1 - c_j$ is reflected specularly, i.e.,

$$f_j(s) = c_j f_j(s) + (1 - c_j) f_j(s - 2(\Delta n) n), \quad \text{vn} > 0,$$

$$f_j(s) = \frac{m_j}{2\pi k_b T}^{\frac{3}{2}} \exp\left(\frac{m_j u_j^2}{2 k_b T}\right), \quad j = m, n,$$

where $n$ is the inner normal to the surface of the capillary; the superscripts $+ \rightarrow -$, $- \rightarrow +$ refer, respectively, to particles that are reflected, emitted diffusively from, and incident upon the surface; $n_j$ is the number density of particles emitted diffusively in the $j$-th state, $m_0$ is the mass of a particle, and $k_b$ is Boltzmann's constant.

We now consider the case of low radiation intensity $I$, in which $\kappa(v) < 1$. Then the states of the components of the gas mixture are slightly nonequilibrium, and the distribution functions $f_m$ and $f_n$ can be represented in the form

$$f_j(s) = f_{j0} [1 + h_j(s, v)],$$

where

$$f_{j0} = n_{j0} \left(\frac{m_0}{2\pi k_b T}\right)^{\frac{3}{2}} \exp\left(\frac{-m_j u_j^2}{2 k_b T}\right), \quad j = m, n,$$

and

$$f_{j0} = \frac{n_{j0}}{2}, \quad j = m, n.$$
The second term on the right-hand side of (7) does not depend on the molecular velocities, and consequently it does not contribute to the macroscopic velocity and stress tensor (6). It will therefore be dropped in the subsequent discussion.

It should be noted that the probability of stimulated transitions is small under low-intensity irradiation. Then the ratio \( n_{ml}/n_{nn} \) is also a small quantity, i.e., \( n_{ml}/n_{nn} \leq 1 \).

Equations (5)–(7) make it possible to determine the functions \( h_j(r,c) \) uniquely. However, it is of practical interest to determine the resulting particle flux (light-induced drift) averaged over the transverse cross section of the capillary:

\[
J = J_m + J_n = 2 \int_0^1 (n_{w,nn} + n_{w,wn}) r dr.
\]

For numerical calculations it is convenient to use the dimensionless quantity \( G \), which is related to the flux \( J \) by

\[
J = \frac{n_R eta}{2} \kappa G,
\]

where

\[
\kappa = \frac{1}{\sqrt{\pi}} \int_0^\infty c_j \exp(-c_j^2) c_j dc_j.
\]

We now transform to a new coordinate system in velocity space: we replace the independent variables \( c, c \) by the variables \( c, \theta \) defined as follows:

\[
c_j = \sqrt{c_j^2 + c_p^2}, \quad \theta = \arctan(c_p/c_j).
\]

Since the longitudinal component \( c_j \) of the molecular velocity appears in Eq. (5) as a parameter, we introduce the following truncated perturbation functions:

\[
\tilde{\Phi}_j(r,c_1, \theta) = \frac{\alpha_j}{\sqrt{\pi}} \int_0^\infty h_j(r,c) \exp(-c_j^2) c_j dc_j,
\]

where \( \alpha_m = 1 \) and \( \alpha_n = n_{nn}/n_{nn} \).

We also introduce, in accordance with (6) and (12), new functions for the macroscopic velocities and viscous stresses:

\[
w_i(r) = \frac{\alpha_j \mu(r)}{\sqrt{\pi}} \int_0^\infty \Phi_i(r,c_1, \theta) \exp(-c_j^2) c_j dc_j d\theta,
\]

\[
t_i(r) = \frac{\alpha_j \mu(r)}{\sqrt{\pi}} \int_0^\infty \Phi_i(r,c_1, \theta) \exp(-c_j^2) c_j^2 dc_j d\theta.
\]

Neglecting terms of order \( \alpha_n \) and taking (11)–(14) into consideration, we obtain from (5) the following kinetic equations for the truncated perturbation functions \( \Phi_{j1} \) and \( \Phi_{j2} \):

\[
c_1 \cos \theta \frac{\partial \Phi_{j1}}{\partial r} - c_1 \sin \theta \frac{\partial \Phi_{j1}}{\partial \theta} + R \Phi_{j1} = -\frac{1}{2} \Gamma_{nn} R \Phi_{j1} + R \Phi_{j1} \Phi_{j1} + R \Phi_{j1} \Phi_{j2} + 2c_1 \circ \Phi_{j1} \cos \theta + 2 R c_1 (1 - \Phi_{j1} + \Phi_{j2} \cos \theta),
\]

\[
c_1 \cos \theta \frac{\partial \Phi_{j2}}{\partial r} - c_1 \sin \theta \frac{\partial \Phi_{j2}}{\partial \theta} + R \Phi_{j2} = -\frac{1}{2} \Gamma_{nn} R \Phi_{j2} + R \Phi_{j2} \Phi_{j2} + 2c_1 \circ \Phi_{j2} \cos \theta + 2 R c_1 (1 - \Phi_{j1} + \Phi_{j2} \cos \theta).
\]

(15)

(16)

On the basis of (12), the boundary conditions (7) assume the form

\[
\Phi_j(r_0, c_1, \theta) = (1 - e_j) \Phi_j(r_0, c_1, \theta), \quad j = m, n.
\]

It is well known that for gas motion in a capillary, the accommodation coefficients \( e_j \) are close to unity, while the relative difference between the effective diameters of excited \( \sigma_m \) and unexcited \( \sigma_n \) particles is small:

\[
1 - e_j < 1, \quad |\Delta \sigma/\sigma_m| < 1, \quad \Delta \sigma = \sigma_m - \sigma_n, \quad j = m, n.
\]

We choose the effective collision rate by analogy with the Bhatnagar–Gross–Krook model in the form \( \gamma_n = p' \theta \), where \( \theta \) is the dynamic viscosity coefficient and \( p \) is the pressure. We model the gas particles as hard elastic spheres with effective diameters \( \sigma_n \) and \( \sigma_m \) for the excited and unexcited particles, respectively. Then the rarefaction parameter \( R \) is related to the Knudsen number by

\[
R = R_n = \frac{\sqrt{\pi}}{2} K n^{-1}, \quad K n = \frac{L}{R_n}.
\]

(19)

By virtue of the inequalities (18), we have

\[
S = \frac{R_n}{R_m} = \frac{\gamma_m}{\gamma_n} = \frac{\sigma_m^2}{\sigma_n^2} > 1, \quad \Delta \sigma = \sigma_m - \sigma_n = \frac{\sigma_m + \sigma_n}{2}.
\]

(20)

As a result, we obtain from Eqs. (15) and (16) for the hard-sphere molecular model

\[
c_1 \cos \theta \frac{\partial \Phi_{j1}}{\partial r} - c_1 \sin \theta \frac{\partial \Phi_{j1}}{\partial \theta} + R \Phi_{j1} = -\frac{1}{2} \Gamma_{nn} R \Phi_{j1} + R \left( w_w \left( \frac{5}{6} \Delta w_w \right) + \frac{1}{2} R c_1 S_t \cos \theta \right),
\]

\[
c_1 \cos \theta \frac{\partial \Phi_{j2}}{\partial r} - c_1 \sin \theta \frac{\partial \Phi_{j2}}{\partial \theta} + R \Phi_{j2} = \frac{1}{2} - \Gamma_{nn} R \Phi_{j2} + \left( \frac{5}{6} S w_w R - \frac{1}{3} R c_1 S_t \cos \theta \right).
\]

(21)

(22)

Here the fact that \( t_w = -w_w \) as a result of momentum conservation has been taken into account.
Linearizing with respect to the small parameters (18) makes it possible to separate the surface and collisions mechanisms of light-induced drift. Then

\[ w_j(r) = w_j(0) \Delta \sigma + w_{j+1}(r) \Delta \sigma, \quad \Delta \sigma = e_{i+1} - e_i, \quad j = 1, m. \]

(23)

\[ G(l, l_0) = G(l, l_0) \Delta \sigma + G_{l+1}(l, l_0) \Delta \sigma. \]

(24)

The kinetic coefficients \( G_1 \) and \( G_2 \), which characterize the respective contributions of the surface and collisional mechanisms to light-induced drift, depend only on the rarefaction parameter \( R \) and the \( \Gamma_{l_0} \), the ratios of the radiative decay rate \( \Gamma_{l_0} \) of the excited level to the intermolecular collision rate \( \gamma_x \). In the approximation (18), the kinetic equations do not depend on \( e_i \) and \( \Delta \sigma / \sigma_i \).

Substituting (24) into Eq. (9) gives the following expression for the light-induced drift flux:

\[ J = \frac{n_0 \Delta \sigma}{2} \left( G_1 \Delta \sigma + G_2 \Delta \sigma \right). \]

(25)

3. SOLUTION OF THE KINETIC EQUATIONS

To solve Eqs. (21) and (22) with the boundary conditions (17), we employ the numerical method of discrete ordinates, which is based on replacing the continuous phase space by a discrete phase space. In so doing it is assumed that the gas molecules cannot move with arbitrary velocities, but only with velocities from a fixed set, which comprises \( c_i; (i = 1, \ldots, N) \). This set forms the nodes of a computational mesh in molecular velocity space.

We also introduce a discrete configuration space with nodes of a computational mesh in molecular velocity space.

Using the zeroth approximation \( p = 0 \) for certain (generally arbitrary) profiles of the partial macroscopic velocities and the stress tensor, the values of \( \Phi_{l, q}^{i+1/e} \) can be obtained by iteration to any higher approximation. The new values of the macroscopic quantities \( w_j^{i+1/e} \) and \( c_i^{i+1/e} \) are calculated from the formulas (28) and (29) at each iteration step \( p > 0 \). The rate of convergence of the iteration process depends largely on the choice of the zeroth approximation. In the present work the equilibrium state of the gas, in which

\[ w_j^{0} = 0, \quad c_i^{0} = 0, \]

was chosen as the zeroth approximation. The iterations were continued until the relative difference between the computed values of the macroscopic parameters in the \( p \)-th and \( (p-1) \)-th approximations was less than a prescribed value \( \epsilon = 10^{-9} \).

A nonuniform mesh, consisting of \( N_j = 11 \) Gaussian nodes, was used for the variable \( c_i \), and a uniform mesh with \( N_j = 20 \) nodes was used for the variable \( \theta \). The mesh in configuration space consisted of \( N_j = 200 \) nodes, distributed uniformly along the radius of the capillary. The computational error was at most 0.1% for any values of the rarefaction parameter \( R \) and rate parameter \( \Gamma_{l_0} \).

The computational results are displayed in Figs. 1–4. Some of the computed values of the kinetic coefficients \( G_1(R, \Gamma_{l_0}) \) and \( G_2(R, \Gamma_{l_0}) \) are presented in Tables I and II.
4. RESULTS AND COMPARISON WITH EXPERIMENT

The direction of the surface component of light-induced drift is determined by the signs of the difference $\Delta \varepsilon = \varepsilon_e - \varepsilon_a$, between the accommodation coefficients of the unexcited and excited particles, and the frequency offset $\Omega = \omega - \omega_a$ of the radiation from the center of the absorption line. If $\Delta \varepsilon > 0$, the surface component of drift is oriented in the direction of the radiation for $\Omega > 0$, and opposite the radiation for $\Omega < 0$.

Figure 1 displays the kinetic coefficient $G_1$, which characterizes the surface light-induced drift velocity as a function of the rarefaction parameter $R$ for various values of the frequency parameter $\Gamma_{\text{max}}$. Clearly $G_1$ decreases monotonically as the Knudsen regime passes into the hydrodynamic regime.

The decrease in surface drift velocity with increasing $\Gamma_{\text{max}}$ and fixed $R$ is explained by a decrease in the relative number of excited particles that collide with the capillary walls.

For fixed $\Gamma_{\text{max}}$, the $R$-dependence of the kinetic coefficient $G_2$, which characterizes the collisional light-induced drift, is nonmonotonic (Fig. 2). The curve $G_2(R)$ has a maximum and a minimum, and for some value $R^{(0)}(\Gamma_{\text{max}})$ of the rarefaction parameter, the coefficient $G_2$ changes sign. This means that the direction of the collisional component of light-induced drift is determined not only by the sign of the frequency offset $\Omega$ from the center of the absorption line, but by the gas pressure in the capillary as well. A possible reason for this behavior of $G_2(R)$ is discussed in detail in Ref. 9 for $\Gamma_{\text{max}} < 1$, where the inversion value $R^{(0)}$ of the rarefaction parameter $R$ is fixed. In the general case, the collisional
TABLE I. Kinetic coefficient $G_2(R,T,\Gamma)$ for surface light-induced drift.

<table>
<thead>
<tr>
<th>$R$</th>
<th>0.01</th>
<th>1</th>
<th>5</th>
<th>7</th>
<th>10</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.01</td>
<td>2.975</td>
<td>2.879</td>
<td>2.585</td>
<td>2.469</td>
<td>2.319</td>
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<tr>
<td>0.04</td>
<td>2.635</td>
<td>2.413</td>
<td>1.834</td>
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<td>1.417</td>
</tr>
<tr>
<td>0.07</td>
<td>2.416</td>
<td>2.124</td>
<td>1.442</td>
<td>1.240</td>
<td>1.020</td>
</tr>
<tr>
<td>0.1</td>
<td>2.249</td>
<td>1.911</td>
<td>1.388</td>
<td>0.992</td>
<td>0.791</td>
</tr>
<tr>
<td>0.2</td>
<td>1.867</td>
<td>1.451</td>
<td>0.738</td>
<td>0.583</td>
<td>0.440</td>
</tr>
<tr>
<td>0.3</td>
<td>1.616</td>
<td>1.174</td>
<td>0.527</td>
<td>0.407</td>
<td>0.300</td>
</tr>
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<td>0.4</td>
<td>1.430</td>
<td>0.983</td>
<td>0.406</td>
<td>0.310</td>
<td>0.227</td>
</tr>
<tr>
<td>0.7</td>
<td>1.068</td>
<td>0.654</td>
<td>0.238</td>
<td>0.179</td>
<td>0.130</td>
</tr>
<tr>
<td>1.0</td>
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<td>0.485</td>
<td>0.168</td>
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<td>0.0912</td>
</tr>
<tr>
<td>2.0</td>
<td>0.494</td>
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<td>0.0840</td>
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<td>0.0456</td>
</tr>
<tr>
<td>3.0</td>
<td>0.343</td>
<td>0.171</td>
<td>0.0560</td>
<td>0.0419</td>
<td>0.0304</td>
</tr>
<tr>
<td>5.0</td>
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<td>0.103</td>
<td>0.0336</td>
<td>0.0251</td>
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<tr>
<td>7.0</td>
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<td>0.073</td>
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</tr>
<tr>
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<td>0.0514</td>
<td>0.0167</td>
<td>0.0125</td>
<td>0.0091</td>
</tr>
</tbody>
</table>

light-induced drift reverses direction at $R^{(0.5)}=0.05-0.25$, depending on the value of the frequency parameter $\Gamma_m$. Thus, the direction of collisional light-induced drift at fixed gas pressure in the capillary is also determined by the magnitude of the radiative decay constant $\Gamma_m$.

It can be seen from Fig. 2 that the greater the value of $\Gamma_m$, i.e., the fewer the collisions an excited particle undergoes on average before being quenched, the smaller the value of the kinetic coefficient $G_2$. In the limit $\Gamma_m=\infty$, there is enough time for all excited particles to decay to the ground state during the time taken to cover one mean free path length, and there is therefore no collisional light-induced drift.

Note that for $\Gamma_m<0.1$, the kinetic coefficients $G_1$ and $G_2$ do not depend on $\Gamma_m$. In this range of values of $\Gamma_m$, the results of Ref. 9 are valid.

Figures 3 and 4 display the dimensionless macroscopic velocities $w_1$ and $w_2$ as functions of the radial coordinate $r$. One can see (Fig. 3) that the surface component $w_1$ of the velocity for $R>1$ remains essentially constant across the capillary. This suggests that the viscosity of the gas does not greatly affect the surface light-induced drift.

Note that the surface drift velocity calculated in Ref. 4 has the typical parabolic profile of standard Poiseuille flow. The evolution of the profile of the collisional component $w_2$ of the velocity as the free-molecular regime passes into the hydrodynamic regime is interesting (Fig. 4). As long as the rarefaction parameter $R$ is small, $w_2$ will be virtually independent of $r$. For $R>0.1$, the structure of the light-induced drift flow becomes more complicated. A “core” of the flow, directed in one direction, forms near the axis of the capillary, while in the layer near the wall the gas moves in the opposite direction. The existence of a counterflow is related to the fact that the flux of the excited particles near the wall for intermediate values of $R$ is higher than the flux of unexcited particles, and conversely near the axis of the capillary. In the almost free-molecular regime ($R\ll0.1$), the velocity $w_2$ is directed opposite the wave vector $k$, and $w_2$ is

TABLE II. Kinetic coefficient $G_2(R,\Gamma_m)$, $10^{-3}$, for collisional light-induced drift.

<table>
<thead>
<tr>
<th>$R$</th>
<th>0.01</th>
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<th>5</th>
<th>7</th>
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<td>0.01</td>
<td>$-1.916$</td>
<td>$-1.581$</td>
<td>$-0.946$</td>
<td>$-0.785$</td>
<td>$-0.619$</td>
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<td>0.03</td>
<td>$-2.995$</td>
<td>$-2.135$</td>
<td>$-0.881$</td>
<td>$-0.620$</td>
<td>$-0.377$</td>
</tr>
<tr>
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<td>$-3.206$</td>
<td>$-2.059$</td>
<td>$-0.558$</td>
<td>$-0.295$</td>
<td>$-0.079$</td>
</tr>
<tr>
<td>0.07</td>
<td>$-3.122$</td>
<td>$-1.787$</td>
<td>$-0.228$</td>
<td>$-0.003$</td>
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<td>$-1.251$</td>
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<td>0.336</td>
<td>0.393</td>
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<td>0.878</td>
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</tr>
<tr>
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<td>0.675</td>
<td>0.396</td>
<td>0.217</td>
</tr>
<tr>
<td>3.0</td>
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<td>0.169</td>
<td>0.091</td>
</tr>
<tr>
<td>7.0</td>
<td>6.254</td>
<td>1.743</td>
<td>0.211</td>
<td>0.120</td>
<td>$-$</td>
</tr>
<tr>
<td>10.0</td>
<td>4.619</td>
<td>1.237</td>
<td>0.147</td>
<td>0.083</td>
<td>0.044</td>
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greater near the wall than on the axis of the capillary (concave profile). In the hydrodynamic regime, for \( R > 10 \), the velocity \( w_2 \) is oriented in the same as the direction as the wave vector \( k \), and \( w_2 \) increases away from the wall.

Light-induced drift has been studied experimentally in CHF\(_3\) molecules. Resonant CO\(_2\) laser radiation was directed along a quartz capillary with radius \( R_0 = 0.75 \) mm and length \( L = 300 \) mm. The mean gas pressure in the cell was varied over the range 0.7–280 Pa, which corresponds to rarefaction parameters \( R = 0.1-47 \).

The results of the measurements are presented in Ref. 13 for the quantity

\[
\frac{\Delta p/P}{\Delta \epsilon/\epsilon} = L \frac{\psi(\Omega)}{2} \left( R_0 + R \right) \left( P_1 \Delta \epsilon + P_2 \Delta \epsilon \right),
\]

where \( \psi(\Omega) \) is an antisymmetric function of the detuning \( \Omega \) (Ref. 14), and \( G_\Omega \) is a kinetic coefficient that characterizes Poiseuille flow and depends on the parameters \( R \) and \( \epsilon_\nu \).

In Fig. 5 the theoretical results are compared with the experiment of Ref. 13. The values chosen for the parameters \( \Delta \epsilon \) and \( \Delta \epsilon/\epsilon \) are presented in Table III. The small discrepancy with the results of Ref. 9 results from our more accurate numerical calculation, and the uncertainty in the values is due to the experimental uncertainty in the homogeneous half-width \( \Gamma \) of the absorption line, and thus the function \( \psi(\Omega) \). The theoretical curves corresponding to the Eq. (31) satisfactorily describe the experimental data at all pressures.

Financial support for this work was provided by the International Science Foundation (grant No. RG4000).

Translated by M. E. Alferieff

<table>
<thead>
<tr>
<th>Branch</th>
<th>( R(4, 3) )</th>
<th>( R(31, 9) )</th>
<th>( P(24, 13) )</th>
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</thead>
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<tr>
<td>( \Delta \epsilon \times 10^{-3} )</td>
<td>-1.65</td>
<td>-2.5-3</td>
<td>-2.1-3.3</td>
</tr>
<tr>
<td>( \Delta \epsilon/\epsilon \times 10^{-3} )</td>
<td>5.08</td>
<td>-50-64</td>
<td>-59.9-90.5</td>
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</tbody>
</table>

FIG. 5. Comparison of theory (solid lines) with experimental data for CHF\(_3\) molecules: 1) \( P(24, 13) \); 2) \( Q(12, 3) \); 3) \( R(4, 3) \); 4) \( R(31, 9) \).