Light-induced particle drift in the presence of quasiresonance energy transfer

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The effect of quasiresonance energy transfer on light-induced particle drift is investigated. Owing to the large cross section ($\lesssim 10^{-12}$ cm^2) for quasiresonance collisions, these collisions slow down the excited-particle flux produced by the field. Nonetheless the magnitude of the drift velocity can be as high as half the mean thermal velocity.

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1. INTRODUCTION

It is shown in Refs. 1 and 2 that a running light wave induces a directed flux of absorbing particles mixed with a buffer gas. The physics of the phenomenon of light-induced drift (LID) consists in the following. Because of the Doppler effect, the excitation of the particles occurs selectively in velocity space, as a result of which there arise oppositely directed fluxes of excited ($\mathbf{j}_m$) and unexcited ($\mathbf{j}_n$) particles. In the case in which the cross sections for elastic collisions of the particles in the ground and excited states are different, the buffer gas exerts different frictional forces on the fluxes $\mathbf{j}_m$ and $\mathbf{j}_n$. As a result, the absorbing gas begins to move as a whole. This phenomenon has been experimentally confirmed for the Na atom and the SF_6 molecule.[4]

In the present paper we investigate the effect of quasiresonant energy transfer on the particle drift. This process is of interest because of the fact that the quasiresonant exchange is characterized by a large cross section[5,6]

$$\sigma_{ex}\sim 10^{-12}-10^{-14}\,\text{cm}^2, \tag{1.1}$$

which significantly exceeds the gaskinetic cross section ($\sim 10^{-15}$ cm^2). The magnitude of the cross section σ_{ex} is largely determined by two factors: the oscillator strengths of the corresponding transitions of the colliding particles and the Massey parameter $\chi=\Delta Ea/\hbar v$. Here ΔE is the resonance defect; a is the critical impact parameter, i.e., the parameter for which the probability for excitation transfer is of the order of unity; v is the relative velocity of the colliding particles. But it is well known from the theory of the elastic LID effect[1,2] that collisions involving a change in the velocity direction are necessary for the appearance of drift. This means that the drift is determined not by σ_{ex}, but by the transport cross section

$$\sigma_{ex}^{t}=\int d\Omega\,(1-\cos\theta)\,\sigma_{ex}(\theta), \tag{1.2}$$

which is naturally smaller than σ_{ex}, since in the case of quasiresonance exchange the scattering is largely small-angle scattering.[7] Let us estimate σ_{ex}^{t} for the dipole–dipole interaction

$$V(r)=C/r^3.$$

The classical scattering angle for this potential depends on the impact parameter ρ as follows[8]:

$$\theta_{cl}\sim(\rho_0/\rho)^3,\quad \rho_0=a(\hbar/\mu va)^{1/4}. \tag{1.3}$$

It can be seen from (1.2) and (1.3) that only flights with $\rho\lesssim\rho_0$ make a contribution to σ_{ex}^{t}. Therefore,

$$\sigma_{ex}^{t}\sim\pi\rho_0^2\sim\sigma_{ex}(\hbar/\mu va)^{1/2}, \tag{1.4}$$

where $\sigma_{ex}=\pi a^2$ and μ is the reduced mass of the colliding particles. For example, for $\mu=20$, $v=10^5$ cm/sec, and $\sigma_{ex}=10^{-12}-10^{-14}$ cm^2 we find from (1.4) that $\sigma_{ex}/\sigma_{ex}^{t}\sim 100-25$, or $\sigma_{ex}^{t}\sim 10^{-14}-0.4\times10^{-15}$ cm^2, which is comparable to, or an order of magnitude greater than the gas-kinetic cross section.

Let us qualitatively consider the drift motion due to the quasi-resonance excitation exchange between the absorbing particles A and the buffer-gas particles B:

$$A^*+B\to A+B^*+\Delta E,$$

where A^* and B^* are excited particles. As in Refs. 1 and 2, here we shall consider the case in which the

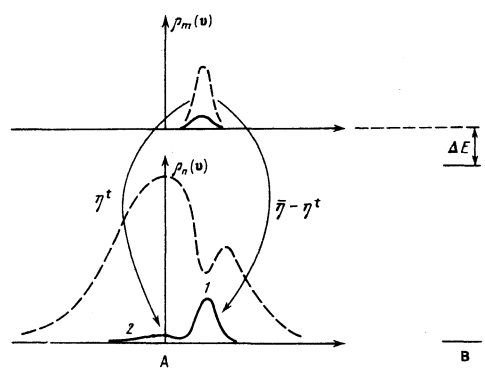

FIG. 1. A denotes the absorbing gas; B, the buffer gas. The dashed curves are plots of the distributions $\rho_m(\mathbf{v})$ and $\rho_n(\mathbf{v})$ without allowance for the collisions with the buffer gas. The solid curves are plots of the distributions $\rho_m(\mathbf{v})$ and $\rho_n(\mathbf{v})$ with allowance for the collisions involving quasiresonance energy transfer (the inelastic channel). 1) The Bennett peak "transferred" to the level n by quasiresonance collisions involving no change in the velocity direction; 2) the particle distribution function transferred to the level n from the level m by quasiresonance collisions involving a change in the velocity direction. The case in which $\gamma<k\bar{v}$ is depicted here.

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buffer-gas density N_B is high compared to the absorbing-gas density N ($N_B \gg N$). This allows us to neglect the inverse reaction

$$A+B^* \to A^*+B-\Delta E.$$

Figure 1 shows a typical excitation-exchange scheme for A^* and B. When the two-level system is acted upon by monochromatic radiation with wave vector $\mathbf{k}$ and frequency ω, a dip is burnt out (at $\mathbf{k}\cdot\mathbf{v}=\Omega\equiv\omega-\omega_{mn}$) in the velocity distribution, $\rho_n(\mathbf{v})$, for the particles at the lower level n, and a Bennett peak appears in the velocity distribution, $\rho_m(\mathbf{v})$, for the particles at the level m (see Fig. 1). In the case of a nonzero detuning Ω, the Bennett peak provides a flux $\mathbf{j}_m$ of particles in the state m. The resonance excitation exchange quenches the upper level with frequency η. Because of the departure of particles from the level m, the current $\mathbf{j}_m$ decreases by $\bar{\eta}\mathbf{j}_m$ ($\bar{\eta}\sim\eta$) in unit time. Arriving at the lower level in unit time, however, is not $\bar{\eta}\mathbf{j}_m$, but the lower current $(\bar{\eta}-\eta^t)\mathbf{j}_m$ (the distribution 1 in Fig. 1). The decrease of the arrival current is due to the fact that the resonance excitation exchange is accompanied by a partial change in the initial direction of motion of the particle (the distribution 2 in Fig. 1). The collision-induced change in the velocity direction is characterized by the transport cross section σ^t_{ex}, or by the collision rate η^t corresponding to this cross section [see (2.7)]. Summing the current leaving the level m as a result of the collisions and the current arriving at the level n, we find that the change per unit time in the total absorbing-particle current $\mathbf{J}$ is equal to $-\eta^t\mathbf{j}_m$. By its physical meaning, this is none other than the frictional force (per unit mass) acting on the particle A during the quasiresonance collisions: $\mathbf{F}_{ex}=-\eta^t\mathbf{j}_m$. Besides the frictional force due to the inelastic quasiresonance-collision channel, there is also the frictional force due to the elastic collisions with the buffer gas: $\mathbf{F}_{el}=-\nu_n\mathbf{J}-(\nu_m-\nu_n)\mathbf{j}_m$ (Refs. 1 and 2). Here ν_i is the frequency of elastic collisions of the particle A in the state i with the particles B,[2] $\mathbf{J}=\mathbf{j}_m+\mathbf{j}_n$. Thus, the resultant frictional force exerted on the particle A by the buffer gas has the form

$$\mathbf{F}=\mathbf{F}_{ex}+\mathbf{F}_{el}=-\nu_n\mathbf{J}-(\nu_m+\eta^t-\nu_n)\mathbf{j}_m. \tag{1.5}$$

Setting[1,2] $\mathbf{F}=0$, we find the drifting absorbing-particle flux

$$\mathbf{J}=-\left(\frac{\nu_m+\eta^t-\nu_n}{\nu_n}\right)\mathbf{j}_m. \tag{1.6}$$

It can be seen from this relation that, if the LID effect due to the elastic collisions is absent (i.e., if $\nu_m=\nu_n$), then the LID effect due to the inelastic quasiresonance-scattering channel (the elastic quasiresonance-scattering channel is included in ν_m) appears.

2. EQUATIONS OF HYDRODYNAMICS

Let us investigate the interaction of a running light wave with two-level particles for the case shown in Fig. 1. The kinetic equations for this system coincide with the equations considered in Ref. 2. The difference lies in the collision integral, S_i, for a particle in the state i and a particle of the buffer gas. Besides the elastic scattering channel (S_i^{el}), we shall consider the inelastic quasiresonance-scattering channel (S_i^{ex}) as well:

$$S_i=S_i^{el}+S_i^{ex}, \quad i=m, n. \tag{2.1}$$

It is natural that the quasiresonance collisions make a contribution to the elastic-collision integral (S_i^{el}) as well. We shall include this contribution in the elastic-scattering cross section σ_{mm}, or in the corresponding collision rate ν_m. The following relations, which reflect the law of particle conservation, hold for the collision integral (2.1):

$$\langle S_i^{el}\rangle=\langle S_m^{ex}+S_n^{ex}\rangle=0, \tag{2.2}$$

where the angle brackets denote integration over the velocities. Since in this paper we consider the case $N_B \gg N$, S_m^{ex} is the term describing the current leaving the level m for n, while S_n^{ex} is the term describing the current arriving at the level n from m as a result of a quasiresonance collision. The absorbing gas's macroscopic characteristics of interest to us: the densities ρ_m, ρ_n and the fluxes $\mathbf{j}_m$, $\mathbf{j}_n$,

$$\begin{aligned} \rho_i&=\int\rho_{ii}(\mathbf{r},\mathbf{v})\,d\mathbf{v}, \quad \rho=\int\rho_{mn}(\mathbf{r},\mathbf{v})\,d\mathbf{v},\\ \mathbf{j}_i&=\int\mathbf{v}\rho_{ii}(\mathbf{r},\mathbf{v})\,d\mathbf{v}, \quad \mathbf{j}=\int\mathbf{v}\rho_{mn}(\mathbf{r},\mathbf{v})\,d\mathbf{v}, \end{aligned} \tag{2.3}$$

are determined by the equations of hydrodynamics (the equation of continuity and the equation of momentum balance). These equations are derived from the kinetic equation for the density matrix $\rho_{ij}(\mathbf{r},\mathbf{v})$ in much the same way as is done in Ref. 2.

Let us consider the case of homogeneous broadening, i.e., the case in which

$$\gamma \gg k\bar{v}, \tag{2.4}$$

where γ is the absorption-line width; $\bar{v}$ is the mean thermal velocity of the particles A. In this limit the equations of continuity,

$$\begin{aligned} (2\gamma_m+\eta)\rho_m&=-2\,\mathrm{Re}(iG^*\rho),\\ (\gamma-i\Omega)\rho&=iG(N-2\rho_m), \end{aligned} \tag{2.5}$$

and momentum balance,

$$\begin{aligned} (2\gamma_m+\nu_m+\bar{\eta})\mathbf{j}_m+\bar{v}^2\nabla\rho_m/2&=-2\,\mathrm{Re}\,(iG^*\mathbf{j}),\\ \bar{v}^2\nabla N/2&=-(\nu_m+\eta^t-\nu_n)\mathbf{j}_m-\nu_n\mathbf{J},\\ (\gamma_i-i\Omega)\mathbf{j}+\bar{v}^2(i\mathbf{k}+\nabla)\rho/2&=iG(\mathbf{J}-2\mathbf{j}_m), \end{aligned} \tag{2.6}$$

are valid. Here $G=Ed_{mn}/2\hbar$; E is the amplitude of the electric field; $2\gamma_m$ is the spontaneous-decay constant for the level m; and d_{mn} is the matrix element of the dipole moment. In the equations (2.5) we have discarded the small terms $\nabla\mathbf{j}_m$ and $(i\mathbf{k}+\nabla)\mathbf{j}$ (see Ref. 2). The kinetic coefficients η, $\bar{\eta}$, and η^t, which, like the ν_i (Ref. 2), are computed with the aid of the expressions for the collision integrals S_i^{el} and S_i^{ex} (Ref. 9), have the form

$$\begin{aligned} \eta&=\xi\int_0^\infty du u^2 u_1 e^{-\mu u^2/2KT}\sigma_{ex}(u_1,u),\\ \bar{\eta}&=\xi\int_0^\infty du u^2 u_1\left(\frac{\mu}{m_B}+\frac{2}{3}\left(\frac{\mu}{m}\right)^2\frac{u^2}{\bar{v}^2}\right)e^{-\mu u^2/2KT}\sigma_{ex}(u_1,u), \end{aligned} \tag{2.7}$$

$$\eta^{\tau}=N_B\frac{8\pi^2}{3}\frac{\mu}{m}\left(\frac{\mu}{2\pi KT}\right)^{3/2}\int_0^{\infty}du u^2u_1^3e^{-\mu u^2/2KT}\sigma_{ex}^{\tau}(u_1,u),$$

$$u_1=\left(u^2+\frac{2}{\mu}\Delta E\right)^{1/2},\quad \mu=\frac{mm_B}{m+m_B},\quad \xi=4\pi N_B\left(\frac{\mu}{2\pi KT}\right)^{3/2},$$

m and m_B are the masses of the absorbing and buffer particles; and σ_{ex} is the cross section for the inelastic quasiresonance-scattering channel.[5] In the formulas (2.7), for $\Delta E<0$ the lower integration limit is $|2\Delta E/\mu|^{1/2}$. Notice that we neglect in Eqs. (2.5) and (2.6) the temperature shift, which occurs here both as a result of the quasiresonance energy exchange between the absorbing particles and the buffer gas[10] and as a result of the LID effect.[11] The relaxation constants γ and γ_1 in the equations for ρ, (2.5), and j, (2.6), are found from the off-diagonal collision integral S_{mn} (Ref. 9). These constants have the following structure:

$$\gamma=\gamma_m+\gamma^{el}+\gamma^{ex},\quad \gamma_1=\gamma_m+\gamma_1^{el}+\gamma_1^{ex}, \tag{2.8}$$

where $\gamma^{ex}(\gamma_1^{ex})$ and $\gamma^{el}(\gamma_1^{el})$ are the contributions to $\gamma(\gamma_1)$ from the inelastic quasiresonance-collision channel and the elastic collisions. Here we give the expressions for only γ^{el} and γ_1^{el}:

$$\gamma^{el}=\xi\int_0^{\infty}du u^3e^{-\mu u^2/2KT}\left[\frac{2\pi\hbar}{i\mu u}\{f_m(u,0)-f_n^*(u,0)\}-\sigma_{mn}(u)\right. \tag{2.9}$$

$$\gamma_1^{el}=\xi\int_0^{\infty}du u^3e^{-\mu u^2/2KT}\left[\left(\frac{\mu}{m_B}+\frac{2}{3}\left(\frac{\mu}{m}\right)^2\frac{u^2}{\bar v^2}\right)\right.$$
$$\times\left\{\frac{2\pi\hbar}{i\mu u}(f_m(u,0)-f_n^*(u,0))-\sigma_{mn}(u)\right\}+\frac{2}{3}\left(\frac{\mu}{m}\right)^2\frac{u^2}{\bar v^2}\sigma_{mn}^{t}(u)\Big],$$

where $f_i(u,\theta)$ is the amplitude of the elastic scattering through the angle θ of the particle A in the state i by the particle B; $\sigma_{mn}(u,\theta)=f_m(u,\theta)f_n^*(u,\theta)$; $\sigma_{mn}(u)=\int d\Omega\sigma_{mn}(u,\theta)$; and σ_{mn}^{t} is the transport cross section corresponding to the differential cross section $\sigma_{mn}(u,\theta)$. It can be seen from the expressions (2.9) that γ and γ_1 are not equal to each other. For example, for the collisions in which the phase memory is preserved (i.e., for which $f_m=f_n$), it follows directly from (2.9) that

$$\gamma^{el}=0,\quad \gamma_1^{el}\sim\nu_m=\nu_n. \tag{2.10}$$

We shall, to simplify the exposition, assume that $\mathrm{Im}\,\gamma=\mathrm{Im}\,\gamma_1=0$, i.e., we shall include the collision-induced shift in the detuning Ω. Let us only note that the deviation from each other of $\mathrm{Im}\,\gamma$ and $\mathrm{Im}\,\gamma_1$ leads to an asymmetry in the dependence of the drift velocity and the density drop on Ω.

3. THE DRIFT VELOCITY AND THE ABSORBING-PARTICLE DENSITY

Let us find the drift velocity $\mathbf{u}=\mathbf{J}/N$. In the steady-flow regime, the space derivatives in (2.6) vanish, and for the drift velocity of the absorbing particles as a whole we obtain the expression

$$\mathbf{u}=-\bar v\frac{\mathbf{k}}{k}\frac{\eta^t+\nu_m-\nu_n}{\eta^t+\nu_m+\nu_n}\frac{\Omega k\bar v}{\gamma_1^2+\Omega^2}\frac{\varkappa}{1+\varkappa}\frac{\gamma+\gamma}{2\gamma}$$
$$\times\left[\varkappa_1+\frac{2\nu_n}{\eta^t+\nu_m+\nu_n}\frac{2\gamma_m+\nu_m+\bar\eta}{2\gamma_m+\eta}\right]^{-1}, \tag{3.1}$$
$$\varkappa=\frac{4|G|^2}{2\gamma_m+\eta}\frac{\gamma}{\gamma^2+\Omega^2}\qquad \varkappa_1=\frac{4|G|^2}{2\gamma_m+\eta}\frac{\gamma_1}{\gamma_1^2+\Omega^2}.$$

Here, in contrast to the elastic LID effect, the quantity $2\gamma_m+\eta$ enters into the expression for the saturation parameter $\varkappa$, since the upper level now gets quenched both as a result of its spontaneous decay ($2\gamma_m$) and as a result of the quasiresonance exchange (η).

If the absorption cell is closed at both ends, then the particles will, on account of the drift motion, overflow to one end of the cell until the reverse diffusion flux counterbalances the drift, i.e., until the total flux becomes equal to zero ($\mathbf{J}=0$). The absorbing-particle density distribution can easily be found from Eqs. (2.5) and (2.6) after setting $\mathbf{J}=0$ in them:

$$\frac{\nabla N}{N}=-\mathbf{k}(\eta^t+\nu_m-\nu_n)\frac{\varkappa}{1+\varkappa}\frac{\Omega(2\gamma_m+\eta)}{\gamma_1^2+\Omega^2}\frac{\gamma+\gamma_1}{2\gamma}$$
$$\times\left[\nu_m+\bar\eta-\eta-\frac{\varkappa}{2(1+\varkappa)}(\eta^t+\nu_m-\nu_n)\left(1+\frac{2\gamma_m+\eta}{\gamma}\frac{\gamma\gamma_1-\Omega^2}{\gamma_1^2+\Omega^2}\right)\right.$$
$$\left.+(1+\varkappa_1)(2\gamma_m+\eta)\right]^{-1}. \tag{3.2}$$

Naturally, in the absence of quasiresonance exchange (i.e., for $\eta=\bar\eta=\eta^t=0$), and for $\gamma=\gamma_1\gg\gamma_m$, the expressions (3.1) and (3.2) coincide with the expressions for u and $\nabla N/N$ from Ref. 2. As can be seen from (3.1) and (3.2), the LID effect is also possible when $\nu_m=\nu_n$ (in this case the elastic LID effect does not occur[1,2]).

Let us determine which values of the drift velocity u and the dimension,

$$l=|\nabla N/N|^{-1}, \tag{3.3}$$

of the layer in which the absorbing atoms gather we can expect for $\nu_m=\nu_n$ and a small quasiresonance-exchange cross section, i.e., for

$$\nu_m\gg\eta^t. \tag{3.4}$$

Let us consider the collisions in which the phase memory is destroyed and those in which the memory is preserved. In the first case, as can be seen from (2.9), $\gamma\sim\gamma_1$. Let us estimate u and l for $\varkappa\sim1$ and $\gamma\sim k\bar v$. If $2\gamma_m\sim\nu_m$, then from (3.1) and (3.2) we obtain

$$u/\bar v\sim l_0/l\sim\eta^t/12\nu_m\sim\sigma_{ex}^t/12\sigma_{mm}^t. \tag{3.5}$$

Here $l_0=\bar v/\nu_n$ is the mean free path in the absence of radiation. In Ref. 12 Kalyazin and Sazonov investigated particle drift due to the characteristics of the inelastic scattering of a light excited particle by a heavy buffer-gas particle (with total conversion of the excitation energy into kinetic energy for the light particle). The dependence (3.5) agrees with the linear dependence of u and l^{-1} on the ratio of the inelastic-scattering cross section to the elastic-scattering cross section (if by σ_{ex} we understand the cross section for the process considered in Ref. 12). If, on the other hand, $2\gamma\leq\eta$, then we have a quadratic dependence on the parameter[1)] $\sigma_{ex}^t/\sigma_{mm}^t$:

$$\frac{u}{\bar v}\sim\frac{l_0}{l}\sim\frac{\eta^t}{4\nu_m}\frac{\eta}{\nu_m}\sim\left(\frac{\sigma_{ex}^t}{\sigma_{mm}^t}\right)^2\frac{\sigma_{ex}}{4\sigma_{ex}^t}. \tag{3.6}$$

Collisions in which the phase memory is preserved (i.e., for which $f_m = f_n$) manifests the Dicke narrowing effect[13]: $\gamma \sim \gamma_m + \eta$, but $\gamma_1 \sim \nu_m$, (2.10). Formally, the condition of applicability of the homogeneous-broadening limit, (2.4), may be violated when $\gamma \sim \gamma_m + \eta$. But the weaker condition $\gamma_m + N_B \langle v\sigma_{mm} \rangle \gg k\bar{v}$ is required for the validity of the homogeneous-broadening approximation. It follows from (3.1) and (3.2) that, for collisions in which the phase memory is preserved, and for which $\Omega \sim \nu_m$ and $\varkappa \sim 1$.

$$\frac{u}{\bar{v}} \sim \frac{l_0}{l} \sim \frac{\sigma_{ex}^t}{\sigma_{mm}^t} \frac{\sigma_{mm}}{16\sigma_{mm}^t}, \tag{3.7}$$

i.e., u and l^{-1} linearly depend on $\sigma_{ex}^t/\sigma_{mm}^t$.

Let us now investigate the LID effect when the quasiresonance exchange predominates:

$$\eta^t + \nu_m \gg \nu_n, \quad \eta \gg \gamma_m. \tag{3.8}$$

The quantity $\eta^t + \nu_m$ (and not η^t) figures in (3.8). This is due to the fact that the quasiresonance exchange increases the $m \to m$ (but not the $n \to n$ when $N_B \gg N$) elastic-collision cross section as well. Besides (3.8), we shall assume that $\sigma^t \ll \sigma$, i.e., that $\eta^t + \nu_m \ll \eta, \bar{\eta}$. Under the condition (3.8) the absorption line broadens largely as a result of the quasiresonance energy transfer. We shall, for simplicity, assume that $\gamma = \gamma_1$ and $\eta = \bar{\eta}$. In the limit (3.8), the expression, (3.1), for the drift velocity assumes the form

$$\mathbf{u} = -\bar{v} \frac{\mathbf{k}}{k} \frac{\Omega k\bar{v}}{\gamma^2 + \Omega^2} \left(\frac{\varkappa}{1+\varkappa} \right) \frac{1}{\varkappa + 2\nu_n/(\eta^t + \nu_m)}. \tag{3.9}$$

It can be seen from (3.9) that, for $\Omega = \gamma = k\bar{v}$, $u \geq \bar{v}/4$ in the range $\alpha \leq \varkappa \leq 1$ ($\alpha = 2\nu_n/(\eta^t + \nu_m) \ll 1$). The drift velocity then assumes its maximum value $\bar{v}/2$ for $\varkappa = \alpha$. This value of the drift velocity exceeds the u value for the elastic LID effect.[1,2] The absorbing-particle density drop, (3.2), is, in the limit (3.8), given by the relation ($l_0 = \bar{v}/\nu_n$):

$$\frac{\nabla N}{N} = -\frac{1}{l_0} \frac{\varkappa}{(1+\varkappa)^2} \frac{\eta^t + \nu_m}{\nu_n} \frac{\Omega k\bar{v}}{\gamma^2 + \Omega^2}. \tag{3.10}$$

For $\Omega = \gamma = k\bar{v}$ and $\varkappa = 1$, we find from (3.10) that

$$\frac{l_0}{l} = \frac{\eta^t + \nu_m}{8\nu_n} \sim \frac{\sigma_{ex}^t}{8\sigma_{nn}^t} \gg 1. \tag{3.11}$$

This relation implies that the absorbing particles can gather in a layer of thickness l smaller than the mean free path (in zero field) l_0. The fact that $l \ll l_0$ does not yet imply the inapplicability of the hydrodynamic description, as furnished by Eqs. (2.5) and (2.6). The point is that the hydrodynamic equations (2.5) and (2.6) are valid for l greater than the mean free path in a medium with a radiation field (l_0'). But in the limit (3.8) l_0' is much shorter than l_0, since the quasiresonance collisions increase the effective frequency of the collisions involving a change in the velocity direction.

CONCLUSION

We have shown, by comparing the elastic LID effect[1,2] with the LID effect due to the quasiresonance collisions, that in the latter case, when $\sigma_{ex}^t \gg \sigma_{nn}^t$, the drift velocity can significantly exceed the drift velocity obtained when only the elastic collisions are taken into consideration, and can be as high as half the mean thermal velocity. The radiation intensity required for this purpose is higher than in the case of the elastic LID effect by a factor of η/γ_m. For example, for $\gamma_m = 10^7$ Hz, $\bar{v} = 10^5$ cm/sec, $\sigma_{ex} = 10^{-13}$ cm^2, and $N_B = 10^{17}$ cm^{-3}, the ratio η/γ_m is equal to 100. An example of a system in which the quasiresonance collisions should be the most clearly manifested may be an isotopic mixture. In this case one of the isotopes absorbs the radiation, while the others play the role of a buffer gas. As such a system we can choose, for example, a mixture of atomic ^{6}Li and ^{7}Li isotopes with $\sigma_{ex} \sim 10^{-13}$ cm^2 (the D line). It should be noted that the concentration of the light-absorbing isotope should be such that the absorption path exceeds l, (3.3). Other schemes are possible besides the above-considered one (see Fig. 1). For example, because of the fine structure of the levels m ($m_1 m_2 \ldots$) and n ($n_1 n_2 \ldots$), quasiresonance collisions involving $m_i \to m_j$ and $n_i \to n_j$ excitation energy transfers to the buffer particle are possible. In this case the drift velocity will be proportional to $(\eta_m^t + \nu_m - \eta_n^t - \nu_n)/(\eta_n^t + \nu_n)$. Such a scheme is possible for the vibrational transitions of atoms in molecules in the presence of strong R-R relaxation or for electronic transitions in the presence of V-V relaxation. In conclusion, let us note that quasiresonance excitation transfer plays the same positive role in diffusional suction and ejection.[14]

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Note. The prediction made in the present paper that drift can occur in an isotropic mixture in which the role of the buffer gas is played by one of the isotopes has recently been experimentally confirmed by Panfilov *et al.* [JETP Lett. Fiz. **33**, 80 (1981) in the isotopic mixture $^{12}CH_3F+^{13}CH_3F$. From the experimental data Panfilov *et al.* determined the relative collision frequency difference, and found it to be equal to $\Delta\nu/\nu = 10^{-2}-10^{-3}$. The theoretical estimate, carried out by us on the basis of (1.4) and the assumption that $\Delta\nu \sim \sigma^t_{ex}$, given the same order of magnitude for $\Delta\nu/\nu$.
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Quasiclassical representation for the cross sections for vibrational excitation of molecules by slow electrons

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Approximate (quasiclassical) formulas are derived which allow the computation of the cross section for vibrational excitation of a diatomic molecule in the ground vibrational state in a resonance collision with a slow electron. A numerical comparison of these formulas with Herzenberg's quantum-mechanical calculation for the case of the scattering of an electron of energy ~2.3 eV by molecular nitrogen is carried out. A comparison with experiment shows that the formulas of the present paper are almost as accurate as the results of the exact calculation.

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The vibrational excitation of diatomic molecules by slow electrons has been the subject of many experimental and theoretical investigations (see, for example, Schulz's[1] and Lane's[2] review articles). It is well known that the majority of the experimental data are well accounted for (both qualitatively and quantitatively) within the framework of the idea that the molecular-ion states, which determine the resonance character of the inelastic processes, are quasistable states. It is assumed that there is formed as a result of the interaction of an electron with a molecule in the ground vibrational state n_0 an intermediate negative ion that decays into the final reaction products[1)]:

$$(AB)_{n_0}+e^-(E_{n_0})\to AB^-\to (AB)_n+e^-(E_n) \qquad (1)$$

(the electron energies are indicated in the brackets, and the energy conservation law for "molecule + electron" system relates E_{n_0} and E_n:

$$E=E_{n_0}+\varepsilon_{n_0}=E_n+\varepsilon_n,$$

where ε_n is the energy of the molecule in its nth vibrational state).

If the lifetime of the group AB^- is short compared to the vibration period of the molecule AB, then the energy dependence of all the processes exhibits normal resonance. If, on the other hand, the lifetime of the group is comparable to the period of the molecular vibrations, then the absorption and emission of the electron in the process (1) occur at different internuclear distances, and there appear in the resonance curve oscillations connected with interference phenomena in the motion of the nuclei.

A method for the theoretical investigation of the processes of inelastic resonance electron scattering by diatomic molecules has been proposed by a number of authors.[3-5] The basic equation in this theory is a Schrödinger-type inhomogeneous equation that, after the separation of the angular dependences, assumes the following form:

$$\left(-\frac{1}{2M}\frac{d^2}{dR^2}+W(R)-E\right)\xi(R)=-\zeta_0(R)V(R). \qquad (2)$$

Here M is the reduced mass of the atoms constituting the molecule; R is the internuclear distance; $W(R)$ is the molecular-ion term [it becomes quasistationary at R less than some R_1: $W(R)=U_I(R)+i\Gamma(R)/2$, where $\Gamma(R)>0$ is the width of this term]; $V(R)$ is the autoionization amplitude of the molecular ion: $V(R)=[4\pi^2\Gamma(R)/mk(R)]^{1/2}$, where $k(R)=[2m(U_I(R)-U(R)]^{1/2}$ is the momentum of the electron emitted by the molecule at an internuclear distance of R; U denotes the term of the neutral molecule and m is the electron mass. The function $\xi(R)$, determined by Eq. (2) with natural boundary conditions, is the wave function for the nuclear motion in the molecular ion; $\zeta_0(R)$ is the wave function for the nuclear motion in the initial (i.e., ground) vibrational state of the molecule.

To determine the cross section for the process (1),

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