The radiation force rectification effect and light-induced transport phenomena in a resonant gas

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The possibility is demonstrated of using intense interfering light fields to form a bulk, rectified, induced-light-pressure force, thereby making it possible to efficiently act on transport phenomena of a small resonant admixture in a buffer gas. This force is capable of inducing rotating or fixed spatially periodic structures in a dense gas. The symmetry and other characteristics of such a sharp periodic stratification of the gas are extremely sensitive to the spatial configuration, phases, and parameters of the fields acting on it. © *1995 American Institute of Physics.*

1. INTRODUCTION

References 1–4 theoretically predicted and Refs. 5 and 6 experimentally confirmed the effect of rectification of a gradient resonant-light-pressure force. This effect arises as a result of the motion of a resonant atom in a bichromatic field with complex amplitude **E** of the form

$$\mathbf{E}(\mathbf{r},t)=(\mathbf{E}_0(\mathbf{r})+\mathbf{E}_1(\mathbf{r})e^{-i\Delta_1 t})e^{-i\Delta_0 t}, \tag{1}$$

where the fields $\mathbf{E}_0(\mathbf{r})$ and $\mathbf{E}_1(\mathbf{r})$ have standing wave structure with characteristic spatial period $\lambda\sim 1/k$, and Δ_0 and $\Delta_1+\Delta_0$ are the detunings from the resonant frequency of the atom. The essence of this effect consists in the appearance of a component of the radiation force F_R having the order of magnitude of the gradient force in a monochromatic standing wave, i.e., $\mathbf{dE}_0 k$ (**d** is the matrix element of the dipole moment operator of the atom), but in contrast to the gradient force it is sign-constant over large (macroscopic) spatial scales ($\gg 1/k$) significantly exceeding the wavelength of the light λ. The rectified gradient force (RGF) can be significantly greater than the spontaneous light pressure force and makes it possible to form deep potential wells and produce various forms of rotational motion of the particles. These unusual properties of the RGF come at a price, namely a bound on the velocity of the atom

$$|v|<\frac{\gamma}{2k}, \tag{2}$$

which delineates the region of maximum values of the RGF in velocity space independent of the field amplitudes (here γ is the spontaneous decay rate of the excited state of the atom). Outside this region the RGF falls rapidly proportional to $(\gamma/kv)^2$ (Ref. 3). Inequality (2) means that spontaneous relaxation processes should take place during the time it takes the atom to move a distance of the order of the characteristic dimensions of a spatial inhomogeneity of the field, $\tau\sim\lambda/v$. This important aspect of the physical mechanism of the formation of the RGF is connected with the need to have incoherent mixing of the adiabatic ("dressed") states of the atom in the field (i.e., the eigenstates of the Hamiltonian of intra-atomic motion in the radiation field).[7] For isolated rarefied ensembles of atoms condition (2) turns out to be automatically fulfilled as the evolution in the resonant field proceeds, thanks to the action of the radiative friction force (even if at first it is violated).[8]

However, quite typical and interesting is the situation with a small relative fraction ($\sim\gamma/ks\ll 1$) of copper atoms when their thermal velocity s is determined by interaction with a thermostat and is large in comparison with their characteristic velocity $\gamma/2k$. An example of a physical object of such kind frequently encountered in laser gas-kinetics is a mixture of a buffer gas with a small resonant admixture. We may ask, are any noticeable manifestations of the rectification of the gradient force possible in this case?

The goal of the present paper is to lay a basis for a positive answer to this question. In this paper we develop a theory of the rectification of the radiative force under conditions of strong collisional relaxation of the translational degrees of freedom, when the mean free path (λ_c) of the resonant atoms in the buffer gas is significantly smaller than the radiation wavelength

$$\lambda_c\ll\lambda. \tag{3}$$

Condition (3) is of fundamental significance and makes it possible to overcome the velocity selectivity of the RGF expressed by inequality (2) (which is valid for a collisionless or weakly collisional gas). Indeed, in this case, the displacement of an atom by a distance of the order of the inhomogeneity scale of the field $k^{-1}\sim\lambda$ will take place in the diffusion regime (but not in the free path regime). If the corresponding characteristic diffusion time is larger than (or of the order of) the atomic relaxation time γ^{-1}, or what is the same thing

$$Dk^2\lesssim\gamma, \tag{4}$$

(where $D\sim\lambda_c s^2$ is the diffusion coefficient of the resonant atoms), then the RGF generated by the bichromatic field will reach its maximum values and act on all the resonant particles identically without regard for the velocity of their Brownian motion. We will show that in this case the RGF provides an extremely efficient means of controlling the transport phenomena of resonant atoms mixed with the buffer gas. The spatial structure of the RGF can be potential or vortical or potential-vortical with controllable ratio of

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weights of the potential and vortical components. In this case the most significant result of the action of the RGF is deep spatially periodic stratification of the resonant gaseous mixture and the formation in it of rotating spatially periodic structures whose shape depends on the configuration and phases of the acting fields.

It is very interesting that in the limit of weak saturation the rectification effect appears in the fourth (!) order of the field amplitude, and not in the sixth as in the collisionless situation.[1] This is due to the influence on the interaction of the atom with the radiation field of collisions with phase jerk.

2. KINETIC EQUATIONS

Let us consider the interaction of a linearly polarized field of the form (1) with resonant two-level atoms which constitute a small admixture to the buffer gas, which is transparent to the radiation.

According to the basic concept of rectification of a gradient force[1–4] the detuning of the field $\mathbf{E}_1$ (which for brevity we will call the high-frequency (HF) field) in the bichromatic field (1) is assumed to be large

$$|\Delta_1| \gg |V_1(\mathbf{r})|,\ |V_0(\mathbf{r})|,\ |\Delta_0|,\ ks,\ \gamma_\perp,\ \nu \tag{5}$$

and its role reduces to the formation of a spatially inhomogeneous Stark level shift $\pm|V_1(\mathbf{r})|^2/\Delta_1$. Here $V_\alpha(\mathbf{r})=\mathbf{d}\mathbf{E}_\alpha(\mathbf{r})/\hbar$ are the local Rabi frequencies, $\alpha=0,\ 1$; $\gamma_\perp$ is the homogeneous resonance width:

$$\gamma_\perp=\frac{\gamma}{2}+\Gamma_c.$$

Γ_c is the frequency of collisions with phase jerk of the atomic oscillator, ν is the characteristic frequency of elastic scattering of the resonant atoms in their collisions with the particles of the buffer gas. We will describe the resonant particles by equations for the density matrix in the Wigner representation with allowance for recoil and collisions,[9] which in the quasiclassical limit ($\hbar k \ll ms$), after averaging over the rapid oscillations with frequency Δ_1 (this procedure is described in detail in Refs. 3 and 8), takes the form

$$\begin{aligned}
&\frac{dp}{dt}+[\gamma_\perp-i\Delta(\mathbf{r})]p=-iVq-\frac{\hbar}{2m}\nabla\frac{\partial f}{\partial \mathbf{v}},\\
&\frac{dq}{dt}+\gamma(f+q)=(2iVp^*+\text{c.c.})\\
&\qquad -\frac{\hbar}{2m}\nabla\Delta(\mathbf{r})\frac{\partial f}{\partial \mathbf{v}}+\left[\frac{\delta q}{\delta t}\right]_c,\\
&\frac{df}{dt}+\frac{\hbar}{m}\left(\nabla V^*\frac{\partial p}{\partial \mathbf{v}}+\text{c.c.}\right)\\
&\qquad =-\frac{\hbar}{2m}\nabla\Delta(\mathbf{r})\frac{\partial q}{\partial \mathbf{v}}+\frac{\delta\hat{\gamma}(f+q)}{2}+\left[\frac{\delta f}{\delta t}\right]_c,\\
&\left[\frac{\delta q}{\delta t}\right]_c=\hat{I}_-(f)+\hat{I}_+(q),\quad \left[\frac{\delta f}{\delta t}\right]_c=\hat{I}_+(f)+\hat{I}_-(q),
\end{aligned}\tag{6}$$

where $p(\mathbf{r},\mathbf{v},t)$ is the nondiagonal element of the density matrix (and has the meaning of the density of the distribution of the induced dipole moment in units of d); $q(\mathbf{r},\mathbf{v},t)$ and $f(\mathbf{r},\mathbf{v},t)$ are, respectively, the difference and sum of the Wigner distributions of the particles in the ground and excited states,

$$\Delta(\mathbf{r})=\Delta_0+\frac{2|V_1(\mathbf{r})|^2}{\Delta_1} \tag{7}$$

is the effective, spatially inhomogeneous detuning of the resonance in the averaged system,[1,2] $\hat{I}_\pm=(\hat{I}_2\pm\hat{I}_1)/2$, where $\hat{I}_{2,1}$ is the integral of the elastic collisions of the excited and unexcited atoms, respectively, with the buffer particles, the operator $\delta\hat{\gamma}$ takes account of recoil in the spontaneous emission of photons,[7]

$$\frac{d}{dt}=\frac{\partial}{\partial t}+\mathbf{v}\frac{\partial}{\partial \mathbf{r}},$$

m is the mass of the resonant atom, and any possible collisional frequency shift is included in Δ_0. Here and in what follows, for brevity we will drop the subscript 0 from V_0. Equations (6) assume strong phase jerk of the atomic oscillator and fulfillability of the condition of the collision approximation for the collision integrals, $|\Delta_1|\tau_c \ll 1$ (Ref. 10) (τ_c is the duration of an individual collision act, the influence of the electromagnetic field on which we will thus neglect). We also neglect collisions of the resonant atoms with each other and also inelastic collisions. Spelling out the physical conditions that are of interest to us, we adopt the following hierarchy of characteristic frequencies of the problem:

$$\begin{aligned}
&\frac{ms^2}{2\hbar}\gg|\Delta|\gg\Gamma_c>\nu,\quad \nu\gg ks\gg\gamma\gg\omega_R=\frac{\hbar k^2}{2m},\\
&\frac{|V_1|^2}{\Delta_1}\gg\gamma>\left|\frac{V}{\Delta}\right|^2\Gamma_c.
\end{aligned}\tag{8}$$

The first group of inequalities in (8) (in combination with the others) means that the average kinetic energy of the resonant atoms is significantly greater than the interaction energy of the atom with the field $\sim\hbar V$, and that the effective detuning Δ is much greater than the velocities of all the relaxation processes. The second group of inequalities corresponds to the basic initial assumption (3) (since $\nu\sim s/\lambda_c$) and to typical relations between physical quantities in the optical spectral range. Finally, the last of conditions (8) ensures that induced light pressure effects will predominate over spontaneous light pressure effects ($|V_1|^2/|\Delta_1|\gamma\gg 1$) (Refs. 1–3, 7) and that the regime of not very strong saturation of absorption will obtain ($|V/\Delta|^2\cdot\Gamma_c/\gamma<1$) (which is necessary, as we will see, for maximization of the RGF).

In addition to the above, relations (8) contain within themselves the conditions of adiabatic motion of the atom in the field, given which the induced dipole moment p for the most part follows the motion of the atom.[7,8] For this reason it turns out to be convenient in the treatment of the problem to transform to the basis of adiabatic states[7] (states "dressed" by the field) by eliminating p from the system of equations (6) by means of the following expansion (for $t>\gamma_\perp^{-1}$) in powers of $1/\Delta(\mathbf{r})$ (Ref. 8):

$$p=p^{(0)}+p^{(1)}+\ldots,\quad p^{(0)}=\frac{Vq}{\Delta(\mathbf{r})},$$

$$p^{(1)}=-\frac{i}{\Delta(\mathbf{r})}\left(\frac{\mathrm{d}}{\mathrm{d}t}+\gamma_{\perp}\right)\frac{Vq}{\Delta(\mathbf{r})}-\frac{\hbar i\nabla V}{2m\Delta(\mathbf{r})}\frac{\partial f}{\partial \mathbf{v}},\tag{9}$$

$$p^{(n)}=-\frac{i}{\Delta(\mathbf{r})}\left(\frac{\mathrm{d}}{\mathrm{d}t}+\gamma_{\perp}\right)p^{(n-1)},\quad n\geqslant 2.$$

Applying this procedure (9) and including terms in the expansion (6) up to second order, inclusive, we obtain the following system of kinetic equations for the auxiliary distribution functions $\rho_{\pm}=(f\pm q\hat{\mu})/2$:

$$\frac{\mathrm{d}\rho_{\pm}}{\mathrm{d}t}\pm\frac{\mathbf{F}_g}{m}\frac{\partial\rho_{\pm}}{\partial\mathbf{v}}\pm[\Gamma_{+}(\mathbf{r})\rho_{+}-\Gamma_{-}(\mathbf{r})\rho_{-}]$$

$$=\left[\frac{\delta\rho_{\pm}}{\delta t}\right]_c-(G_a\pm G_d),\tag{10}$$

where

$$\mu(\mathbf{r})=\sqrt{1+\frac{4|V(\mathbf{r})|^2}{\Delta^2(\mathbf{r})}},\quad \hat{\mu}\simeq\mu(\mathbf{r})\left(1+\frac{2\mathbf{J}\mathbf{v}}{\Delta^3(\mathbf{r})}\right),$$

$$\mathbf{J}=i(V(\mathbf{r})\nabla V^*(\mathbf{r})-\text{c.c.}),\quad \mathbf{F}_g=\hbar\,\mathrm{sign}(\Delta)\frac{\nabla\epsilon(\mathbf{r})}{2},$$

$$\epsilon(\mathbf{r})=\sqrt{4|V(\mathbf{r}|^2+\Delta^2(\mathbf{r})}$$

$$\Gamma_{\pm}(\mathbf{r})=\gamma\left[\left(\frac{1\pm\mu(\mathbf{r})}{2\mu(\mathbf{r})}\right)^2+\frac{\Gamma_c}{2\gamma}\frac{\mu^2(\mathbf{r})-1}{\mu^2(\mathbf{r})}\right],\tag{11}$$

$$\left[\frac{\delta\rho_{\pm}}{\delta t}\right]_c=\frac{1}{2}\left[\left(1+\frac{1}{\mu^2(\mathbf{r})}\right)\hat{I}_{+}\pm\frac{2}{\mu(\mathbf{r})}\hat{I}_{-}\right]\rho_{\pm}$$

$$+\frac{1}{2}\left(1-\frac{1}{\mu^2(\mathbf{r})}\right)\hat{I}_{+}\rho_{\mp},$$

$$G_a\simeq 2\hat{I}_{-}(Q\mathbf{J}\mathbf{v})/\Delta^3(\mathbf{r})\mu^2(\mathbf{r}),$$

$$G_d\simeq 4Q\Gamma_c(\mathbf{J}\mathbf{v})/\Delta^3(\mathbf{r})\mu^2(\mathbf{r}),\quad Q=(\rho_{+}-\rho_{-}).$$

Equations (10) generalize the well-known equations of bipotential kinetics of atoms in a monochromatic standing-wave field (see Ref. 11, p. 117 and also Ref. 7, p. 135). The newly introduced functions $\rho_{\pm}(\mathbf{r},\mathbf{v},t)$ differ from the populations of the "true" adiabatic states by small corrections in the adiabaticity parameter $\sim|V/\Delta|^2kv/\Delta\ll 1$. Terms proportional to $G_a\pm G_b$, of like order of magnitude, which take account of the second and third groups of inequalities (8), are present in the kinetic equation. They depend on the scalar product of the vector $\mathbf{J}$ (which, as is not hard to verify for linearly polarized fields, is proportional to the energy flux density of the radiation at the point $\mathbf{r}$) and the velocity $\mathbf{v}$ and consequently describe kinetic phenomena associated with the velocity selectivity of the interaction of the atom with the resonant field due to the Doppler effect. We also see from Eq. (10) that incoherent mixing of the adiabatic states is realized not only thanks to spontaneous relaxation, but also to collisions. This circumstance is a reflection of the well-known fact that elastic collisions of atoms with particles of the buffer gas are inelastic for the combined system "atom+field."[12] Finally we may note that besides the spontaneous light pressure force $\mathbf{F}_s$, we also neglect the light-induced Lorentz force[3,7,8] $F_L\sim|V/\Delta|^2\hbar k^2v$:

$$|F_s/F_g|\sim(|V_1|^2/\Delta_1\gamma)^{-1}\ll 1,$$

$$|F_L/F_g|\sim|V/\Delta|^2ks/\epsilon<\gamma/|\Delta|\ll 1.$$

3. TRANSPORT EQUATIONS

By virtue of conditions (8) and (3) and the kinetic equations (10), those terms dominate which are associated with the collision integrals $[\delta\rho_{\pm}/\delta t]_c$. This means that the collisions are frequent and that at the diffusion stage of the evolution of the resonant admixture ($t>\nu^{-1}$) the distribution functions $\rho_{\pm}(\mathbf{r},\mathbf{v},t)$ are near their equilibrium values, i.e., they allow (according to the well-known approaches of kinetic theory[13]) the following representation:

$$\rho_{\pm}=W_0n_{\pm}(1+\psi_{\pm}),\quad |\psi|\ll 1,\tag{12}$$

where $n_{\pm}$ is the density of atoms in the adiabatic states, W_0 is the Maxwellian velocity distribution

$$W_0(v)=\frac{1}{(\sqrt{\pi}s)^3}\exp[-(v/s)^2],\quad s=\sqrt{\frac{2T}{m}},\tag{13}$$

T is the temperature of the gas in energy units, and the functions $\psi_{\pm}(\mathbf{r},\mathbf{v},t)$ define the nonequilibrium part of the distributions and in the approximation of linearizing in velocity have the form

$$\psi_{\pm}=\frac{\mathbf{J}_{\pm}\mathbf{v}}{Tn_{\pm}},\quad \mathbf{J}_{\pm}(\mathbf{r},t)=\int\mathbf{v}\rho_{\pm}d\mathbf{v}.\tag{14}$$

Note that under the above conditions the diffusion velocities of the particles are always small in comparison with their mean thermal velocity, i.e., $|J_{\pm}|\ll n_{\pm}s$.

Integrating Eqs. (10) over velocity and repeating the given procedure after multiplying Eqs. (10) by $\mathbf{v}$, we easily obtain with the help of expansion (12) a closed system of equations for the macroscopic variables $n_{\pm}$ and $\mathbf{J}_{\pm}$. We write it out in the variables $n_{+}+n_{-}=n$, $\mathbf{J}_{+}+\mathbf{J}_{-}=\mathbf{J}_a$ (the flux density of the resonant atoms) and $n_{+}-n_{-}=N$, $\mathbf{J}_{+}-\mathbf{J}_{-}=\mathbf{J}_d$ (the flux density of the difference of populations of the adiabatic atomic states):

$$\frac{\partial N}{\partial t}+\mathrm{div}\,\mathbf{J}_d+\Gamma(\mathbf{r})N+\Lambda(\mathbf{r})n=-\frac{4\Gamma_c}{\Delta^3(\mathbf{r})\mu^2(\mathbf{r})}(\mathbf{J}\mathbf{J}_d),\tag{15}$$

$$\frac{\partial n}{\partial t}+\mathrm{div}\,\mathbf{J}_a=0,\tag{16}$$

$$\frac{\partial\mathbf{J}_d}{\partial t}+\Gamma_d(\mathbf{r})\mathbf{J}_d+\Gamma_a(\mathbf{r})\mathbf{J}_a=\left(\frac{\mathbf{F}_g}{m}n-\frac{s^2}{2}\nabla N\right)+C_d\mathbf{J},\tag{17}$$

$$\frac{\partial\mathbf{J}_a}{\partial t}+\nu_{+}\mathbf{J}_a+\frac{\nu_{-}}{\mu(\mathbf{r})}\mathbf{J}_d=\left(\frac{\mathbf{F}_g}{m}N-\frac{s^2}{2}\nabla n\right)+C_a\mathbf{J},\tag{18}$$

where (taking Eq. (8) into account: $\nu_{+}\gg\Gamma_c|V/\Delta|^2$)

$$\Gamma_d\simeq\frac{\nu_{+}}{\mu^2(\mathbf{r})},\quad \Gamma_a=\frac{\nu_{-}}{\mu(\mathbf{r})}+\Lambda(\mathbf{r}),\quad \Lambda=\frac{\gamma}{\mu(\mathbf{r})},$$

$$\Gamma = \frac{1}{\mu^2(\mathbf{r})}\left(\gamma + \gamma_\perp \frac{4|V(\mathbf{r})|^2}{\Delta^2(\mathbf{r})}\right)$$

and within the framework of approximation (12), (14) we have introduced the effective frequencies of the collisions with momentum transfer[13]

$$\nu_\pm = \frac{4}{3} n_\beta \sqrt{\frac{8T}{\pi M}} \frac{m_\beta}{m+m_\beta} \int_0^\infty \xi^5 \exp(-\xi^2)\sigma_\pm^t(g)d\xi,$$

$$g = \sqrt{\frac{2T}{M}}\xi, \quad \sigma_\pm^t = 2\pi \int_0^\pi \sigma_\pm(g,\chi) \times (1-\cos\chi)\sin\chi d\chi,$$

$\sigma_\pm = (\sigma_2 \pm \sigma_\perp)/2$ are the half-sum and half-difference of the differential cross sections of elastic scattering of resonant atoms in the ground "1" and excited "2" states for collisions with particles (with mass m_β) of the buffer bas with density n_β, $M = mm_\beta/(m+m_\beta)$,

$$C_a = \frac{s^2 N \nu_-}{\Delta^3(\mathbf{r})\mu^2(\mathbf{r})}, \quad C_d \simeq -\frac{2s^2\Gamma_c N}{\Delta^3(\mathbf{r})\mu^2(\mathbf{r})} \quad \text{for } \Gamma_c \gg \nu_+.$$

Expressions for the stationary and quasistationary ($|\partial \mathbf{J}_{a,d}/\partial t| \ll \nu_+ |\mathbf{J}_{a,d}|$ for $t \gg \nu^{-1}$) the resonant radiation follow from Eqs. (17) and (18):

$$\mathbf{J}_d = \frac{\mu^2(\mathbf{r})}{\bar{\nu}_+}\left[\frac{\mathbf{F}_g}{m}\left(n + \frac{\gamma+\nu_-}{\mu(\mathbf{r})\nu_+}N\right) - \frac{s^2}{2}\left(\nabla N - \frac{\gamma+\nu_-}{\mu(\mathbf{r})\nu_+}\nabla n\right)\right] - \frac{2s^2\Gamma_c N}{\bar{\nu}_+\Delta(\mathbf{r})}\mathbf{J}, \tag{19}$$

$$\mathbf{J}_a = \frac{\mathbf{F}_{\text{eff}}}{m\bar{\nu}_+} n - D\nabla n, \quad \mathbf{F}_{\text{eff}} = \mathbf{F} + \mathbf{F}_1 + \mathbf{F}_2, \tag{20}$$

where

$$\mathbf{F} = \mathbf{F}_g\left(\hat{N} - \frac{\mu(\mathbf{r})\nu_-}{\nu_+}\right), \quad \mathbf{F}_1 = T\mu(\mathbf{r})\left(\frac{\nu_-}{\nu_+ n}\right)\nabla N,$$

$$\mathbf{F}_2 = T\hat{N}\mathbf{J}\frac{4\Gamma_c\nu_-}{\mu(\mathbf{r})\Delta^3(\mathbf{r})\nu_+},$$

and $\bar{\nu}_\pm = \nu_+(1 - \nu_-^2/\nu_+^2 - \nu_-\gamma/\nu_+^2)$, $\hat{N} = N/n$ is the relative difference of populations of the adiabatic states, and $D = s^2/2\bar{\nu}_+$ is the diffusion coefficient.

From Eqs. (20) we see that the resonant atom flux $\mathbf{J}_a$ is generated by three different effective forces $\mathbf{F}$, $\mathbf{F}_1$, and $\mathbf{F}_2$.

The force $\mathbf{F}$ is the induced-light-pressure force, whose order of magnitude ($F \sim \hbar k \epsilon$) is determined by the characteristic rate of the induced transitions between the atomic states.

The effective forces $\mathbf{F}_1$ and $\mathbf{F}_2$ are of a qualitatively different nature, not associated with light pressure phenomena. The force $\mathbf{F}_1$ is due to the difference in the collision frequencies (and consequently in the diffusion coefficients) of the atoms in different adiabatic states and to the spatial modulation of the difference in the populations of these states [$N = N(\mathbf{r})$] thanks to the inhomogeneity of the field. The result of its action is analogous to the effect of light-induced diffusional injection and ejection, which was first considered in Refs. 14 and 15 for the case of a monochromatic light beam. The force $\mathbf{F}_2$ is proportional to the radiation flux density $\mathbf{J}(\mathbf{r})$ and is also associated with the difference in the transport collision frequencies ($\nu_- \neq 0$). The second factor making up this force is the velocity selectivity of excitation of the atom, due in our case of large fields and large detunings to non-adiabatic corrections to the induced dipole moment, which are described by expansion (9). In other words, the force $\mathbf{F}_2$ is responsible for the appearance under the conditions under consideration of the light-induced drift effect.[16] The effective force $\mathbf{F}_{\text{eff}}(\mathbf{r})$, in addition to its direct dependence on the acting fields and their gradients [through the functions $\epsilon(\mathbf{r})$, $\mu(\mathbf{r})$, $\nabla\epsilon(\mathbf{r})$, and $\mathbf{J}(\mathbf{r})$] also contains a dependence on the state of the internal degrees of freedom of the resonant atom through the population difference of the adiabatic states $N = N(\mathbf{r})$. Such a correlation between the internal and translational degrees of freedom in the problem under consideration is taken into account by the full system of coupled equations (15), (16), (19), (20).

4. RECTIFIED FORCES AND FLUXES

We will investigate the action on the resonant admixture of spatially inhomogeneous fields (of the type standing waves), which can be formed by the appropriate superposition of plane waves.[1–3] The intensities of such fields, which are proportional to $|V_1(\mathbf{r})|^2$ and $|V_0(\mathbf{r})|^2$, and the radiation flux density, which is proportional to $\mathbf{J}(\mathbf{r})$, are periodic or quasiperiodic functions oscillating over spatial scales of the order of the wavelength of the light $\lambda \sim k^{-1}$. We will call this spatial scale microscopic. The difference in the frequencies of the waves forming the resonant field, the small misalignment of their wave vectors, the boundedness of the transverse dimensions of real laser beams give rise to the existence of a macroscopic spatial scale $L \gg \lambda$. The presence of two radically different scales can be formally reflected by introducing dimensionless fast ($\mathbf{r}_1$) and slow ($\mathbf{r}_0$) spatial variables

$$\mathbf{r}_1 = k\mathbf{r}, \quad \mathbf{r}_0 = \delta k \mathbf{r}, \quad \frac{\delta k}{k} = \alpha \sim \frac{\lambda}{L} \ll 1. \tag{21}$$

The coefficients of the system of transport equations (15), (16), (19), (20) depend explicitly on both the fast and the slow variables. The expansion of these coefficients in powers of the field contains nonlinear interference terms (proportional, in particular, to products of the form $|V_1(\mathbf{r})|^2 \mathbf{J}(\mathbf{r})$, $|V_0(\mathbf{r})|^2 \nabla |V_1(\mathbf{r})|^2$, etc.), which generate components of the effective forces and fluxes, which depend only on the slow spatial variable, i.e., we have rectified forces and fluxes. Finding these terms would allow one, under the conditions under consideration, to describe the influence of nonlinear interference effects associated with the light pressure[2,3] on the transport phenomena.

Let the relative difference of the transport collision frequencies be small

$$\nu_-/\nu_+ \ll 1 \tag{22}$$

(in many real situations $\nu_-/\nu_+ \lesssim 0.2$, Ref. 17).

Then from conditions (8) and (22) it follows that the characteristic diffusion velocity in the transport of atoms over the microscopic spatial scale is substantially greater than the light-induced drift velocity

$$Dk \gg \frac{|F_{\text{eff}}|}{m\nu_+}. \tag{23}$$

Inequality (23) means that the depth of the potential wells formed by the rapidly oscillating components of the light-induced forces $\tilde{F}_{\text{eff}}$ is small (of the order of the parameters $\hbar\epsilon/T \ll 1$ and $\nu_-/\nu_+ \ll 1$) in comparison with the kinetic energy of the resonant particles. Therefore they can give rise to only an insignificant spatial modulation of the density of atoms against the background of its variation over macroscopic scales. This circumstance allows us to separate the rapidly oscillating and slow components of the density (flux) of atoms and represent the solution of the transport equations in the form of the following expansion (in fact, in the small parameters $\hbar\epsilon/T$, γ/Δ, ν_-/ν_+, and λ/L):

$$\begin{gathered} n=\bar{n}(\mathbf{r}_0)+\tilde{n}^{(1)}(\mathbf{r}_1,\mathbf{r}_0)+\cdots, \quad |\tilde{n}^{(1)}|\ll\bar{n}, \\ \langle\tilde{n}^{(1)}\rangle=0, \quad N=N^{(0)}(\mathbf{r}_1,\mathbf{r}_0)+N^{(1)}(\mathbf{r}_1,\mathbf{r}_0)+\cdots, \\ N^{(1)}\ll N^{(0)}, \end{gathered} \tag{24}$$

where the angular brackets denote averaging over the microscopic spatial oscillations

$$\langle\cdots\rangle=\lim_{R^3\to\infty}\int(\cdots)\frac{d\mathbf{r}_1}{R^3}, \tag{25}$$

(R^3 is the volume of the integration region over the fast variables).

We will limit ourselves in what follows to the stationary regime of interaction of the radiation with the gas and the case of not-too-strong saturation

$$\frac{|V_1|^2}{\Delta_1\Delta_0}, \quad \frac{|V|^2\Gamma_c}{\Delta_0^2\gamma}\ll 1. \tag{26}$$

Thus, substituting expression (24) in Eqs. (15) and (19), we obtain first-order equations in the spatial distribution of the population difference of the adiabatic states, with allowance for "fast" (small-scale) diffusion

$$\mathbf{J}_d=-Dk\nabla_1 N^{(0)}+\mathbf{J}_d^{(0)}+\cdots, \quad |\mathbf{J}_d^{(0)}|\ll Dk|\nabla_1 N^{(0)}|, \tag{27}$$

$$N^{(0)}=\bar{n}(-1+\hat{N}_1), \quad Dk^2\nabla_1^2\hat{N}_1-\gamma\hat{N}_1=-\Gamma_c\frac{4|V(\mathbf{r})|^2}{\Delta_0^2}. \tag{28}$$

Here and below the subscript "1" ("0") on the differential operators denotes differentiation with respect to only the fast (slow) spatial variables. Similarly, from Eqs. (20), (16), and (24) we obtain equations for the main terms of the expansion of the density and flux of the resonant atoms

$$\begin{gathered} \mathbf{J}_a=\mathbf{J}_R(\mathbf{r}_0)+\tilde{\mathbf{J}}_a(\mathbf{r}_1,\mathbf{r}_0)+\tilde{\mathbf{J}}_a^{(1)}(\mathbf{r}_1,\mathbf{r}_0)+\cdots, \\ |\tilde{\mathbf{J}}_a^{(1)}|\ll|\mathbf{J}_R+\tilde{\mathbf{J}}_a|, \quad \langle\tilde{\mathbf{J}}_a\rangle=0, \end{gathered} \tag{29}$$

$$\operatorname{div}_1\tilde{\mathbf{J}}_a=0, \quad \tilde{\mathbf{J}}_a=\frac{\tilde{\mathbf{F}}_{\text{eff}}^{(0)}}{m\nu_+}\bar{n}-Dk\nabla_1\tilde{n}^{(1)}, \tag{30}$$

$$k\operatorname{div}_1\tilde{\mathbf{J}}_a^{(1)}=-\delta k\operatorname{div}_0\mathbf{J}_R, \tag{31}$$

$$\mathbf{J}_R=\frac{\langle\mathbf{F}_{\text{eff}}^{(0)}\rangle}{m\nu_+}\bar{n}-D\,\delta k\nabla_0\bar{n}, \tag{32}$$

where the tilde above a function symbol indicates its rapidly oscillating part with zero mean, $\mathbf{F}_{\text{eff}}^{(0)}$ is the effective force to the first nonzero order in the perturbation theory sense, and $\langle\mathbf{F}_{\text{eff}}^{(0)}\rangle$ and $\tilde{\mathbf{F}}_{\text{eff}}^{(0)}$ are its smooth (rectified) and rapidly oscillating components. Restricting ourselves to the first nonvanishing terms of the expansion of $\mathbf{F}_{\text{eff}}^{(0)}$ in powers of the field amplitude, we obtain the following expressions for $\tilde{\mathbf{F}}_{\text{eff}}^{(0)}$ and $\langle\mathbf{F}_{\text{eff}}^{(0)}\rangle$:

$$\tilde{\mathbf{F}}_{\text{eff}}^{(0)}=\tilde{\mathbf{F}}_1+\tilde{\mathbf{F}}_2+\tilde{\mathbf{F}}, \quad \tilde{\mathbf{F}}_1=T\frac{\nu_-}{\nu_+}k\nabla_1\hat{N}_1(\mathbf{r}), \tag{33}$$

$$\begin{gathered} \tilde{\mathbf{F}}_2=-T\frac{4\Gamma_c\nu_-}{\Delta_0^3\nu_+}\tilde{\mathbf{J}}, \quad \tilde{\mathbf{F}}=-\hbar k\nabla_1\left(\frac{|V_1(\mathbf{r})|^2}{\Delta_1}+\frac{|V(\mathbf{r})|^2}{\Delta_0}\right), \\ \langle\mathbf{F}_{\text{eff}}^{(0)}\rangle=\mathbf{F}_R+\mathbf{F}_{1R}+\mathbf{F}_{2R}, \end{gathered} \tag{34}$$

$$\mathbf{F}_R=\hbar k\left\langle\hat{N}_1\nabla_1\frac{|V_1(\mathbf{r})|^2}{\Delta_1}\right\rangle, \tag{35}$$

$$\mathbf{F}_{1R}=T\frac{\nu_-k}{\nu_+}\left\langle\frac{2|V(\mathbf{r})|^2}{\Delta_0^2}\nabla_1\hat{N}_1\right\rangle, \tag{36}$$

$$\mathbf{F}_{2R}=\mathbf{F}_{2R}^{(2)}+\mathbf{F}_{2R}^{(4)}, \quad \mathbf{F}_{2R}^{(2)}=-\frac{4T\Gamma_c\nu_-}{\nu_+\Delta_0^3}\langle\mathbf{J}(\mathbf{r})\rangle, \tag{37}$$

$$\mathbf{F}_{2R}^{(4)}=\frac{4T\Gamma_c\nu_-}{\nu_+\Delta_0^3}\left\langle\left(\hat{N}_1+\frac{2|V(\mathbf{r})|^2}{\Delta_0^2}+\frac{6|V_1(\mathbf{r})|^2}{\Delta_1\Delta_0}\right)\mathbf{J}(\mathbf{r})\right\rangle, \tag{38}$$

where $\hat{N}_1$ is found from Eq. (28) and we use the subscript "R" to denote the *rectified forces*. It is easy to see [from Eq. (33)] that the forces $\tilde{\mathbf{F}}$ and $\tilde{\mathbf{F}}_1$ are potential, and the force $\tilde{\mathbf{F}}_2$ (determined by the radiation flux density, which is proportional to $\mathbf{J}$) is vortical ($\operatorname{div}_1\tilde{\mathbf{F}}_2=0$), wherefore the solution of Eq. (30) for the oscillating components of the density and flux has the form

$$\begin{gathered} \tilde{n}^{(1)}=-\bar{n}\frac{\hbar}{T}\left(\frac{\widetilde{|V(\mathbf{r})|^2}}{\Delta_0}+\frac{\widetilde{|V_1(\mathbf{r})|^2}}{\Delta_1}\right)+\frac{\nu_-}{\nu_+}\tilde{\hat{N}}_1\bar{n}, \\ \tilde{\mathbf{J}}_a=\frac{\tilde{\mathbf{F}}_2}{m\nu_+}. \end{gathered} \tag{39}$$

From the condition of solvability of the equation of the second approximation (31) we obtain the equation of continuity of the flux of atoms

$$\operatorname{div}_0\mathbf{J}_R=0. \tag{40}$$

Equations (40), (28), and (32) form a closed system for determining the macrostructure of the density $\bar{n}$ and flux $\mathbf{J}_R$ distributions of the resonant atoms, created by the light-induced rectified forces (35)–(38). We see from the expressions that we have obtained that the rectified gradient force $\mathbf{F}_R$ is determined by the spatial correlator of the relative population $\hat{N}_1(\mathbf{r})$ of the anomalously adiabatic state[7] and the intensity gradient of the HF field $\langle\hat{N}_1\nabla|V_1(\mathbf{r})|^2\rangle$. The spatial dependence of $\hat{N}_1(\mathbf{r})$ is governed by Eq. (28) and through it by the field $V(\mathbf{r})$ independent of $V_1(\mathbf{r})$, from which it is obvious that in the most general case the correlator $\langle\hat{N}_1\nabla|V_1(\mathbf{r})|^2\rangle\neq 0$ and consequently $\mathbf{F}_R$ is also different from zero.

In contrast to the frequencies of the elastic collisions of the excited and unexcited atoms ($\nu_-\neq 0$) expressions (36)–(38) allow for possible contributions to the effective rectified forces $\mathbf{F}_{1R}$ and $\mathbf{F}_{2R}$ due to the effects of light-induced diffusional injection and ejection and light-induced drift. However, it may be easily observed from Eqs. (28) and (36) that the force $\mathbf{F}_1$ is clearly not rectified ($\mathbf{F}_{1R}\rightarrow 0$) independent of the structure of the fields in the limit of frequent collisions, for which the amplified condition (4) ($Dk^2\ll\gamma$) is fulfilled and the population of the adiabatic state follows the field $\hat{N}_1=\hat{N}_1(|V(\mathbf{r})|^2)$.

It can be shown that the analogous assertion regarding $\mathbf{F}_{1R}$ holds, and in the general case in which we do not use the expansion over powers of the field amplitude.

The role of all three components of the effective field in transport phenomena, thus, can depend substantially on the concrete physical conditions and, in particular, as we will see in the following section, on the spatial structure of the field.

5. TRANSPORT PHENOMENA

Let us consider how manifestations of the rectified forces depend on concrete configurations of the fields acting on the gas.

Rectified forces in one dimensional standing waves

Let the fields have the following form:

$$V(\mathbf{r})=V_0\cos(kx),$$

$$V_1(\mathbf{r})=V_1\cos[(k+\delta k)x+\phi],\quad \delta k\ll k. \tag{41}$$

Then from Eqs. (35)–(38) and (28) we easily find

$$F_{1R}=F_{2R}=0,\quad F_R=-\frac{\partial U_R}{\partial x},$$
$$U_R=U_0\cos(2\delta kx+2\phi), \tag{42}$$

where

$$U_0=\frac{\hbar k}{\delta k}\left|\frac{V_0}{\Delta_0}\right|^2\frac{|V_1|^2}{2\Delta_1}\,\frac{\Gamma_c}{4Dk^2+\gamma}.$$

Disappearance of the force F_{2R} is obviously connected with the absence of radiation flux in the field configuration (41): $\mathbf{J}=0$, and disappearance of the force F_{1R}, with the fact that the variable $\hat{N}_1=4\Gamma_c|V(\mathbf{r})|^2/(4Dk^2+\gamma)\Delta_0^2$ "tracks" the field intensity $V(\mathbf{r})$ (in the given case independently (!) of the relation between the parameters Dk^2 and γ).

The depth of modulation of the potential relief induced by the rectified force (with spatial period $\pi/\delta k$) is proportional to the large parameter (the ratio of the macroscopic to the microscopic scale $k/\delta k\gg 1$) and is maximum upon fulfillment of condition (4).

Note that thanks to the presence of collisions with phase jerk ($\Gamma_c\neq 0$) the rectified gradient force arises in the fourth order of the expansion in powers of the field amplitude (but not in the sixth as in the case of a collisionless gas[1]). Therefore, for one and the same field amplitude and $Dk^2<\gamma$ the depth of the potential well (42) is large in comparison with the depth of the macroscopic potential wells in the collisionless situation with the large parameter $(\gamma|V_0|^2/\Delta_0^2\Gamma_c)^{-1}$. The similar effect of the influence of collisions with phase jerk on the radiation force was, in essence, taken into account and discussed already in Ref. (18), where allowance was made for the situation with $\gamma_\perp>\gamma/2$.

Periodic stratification of a resonant gas

The solution of the averaged transport equations (32) and (40) in the case of as field of the form (41) has the form of the Boltzmann distribution

$$\bar{n}=n_0\exp[-U_R(\mathbf{r})/T],\quad J_R=0, \tag{43}$$

where U_R is given by formula (42). The resonant admixture is thus grouped into sharp, well-defined layers periodically arranged along the x axis if the condition $U_0\gtrsim T$ of capture of particles into the potential wells is fulfilled. In this case the characteristic width of a layer $x_0>1/\sqrt{k\delta k}$ and, consequently, is significantly greater than the wavelength of the light, but substantially smaller than the macroscopic period $L=\pi/\delta k$: $\lambda\ll x_0\ll L$. Simultaneously, according to Eq. (39), there exists a small-scale modulation of the density with period of the order of the wavelength of the light

$$\bar{n}^{(1)}=\bar{n}(x)\frac{\hbar}{2T}\left[\frac{|V_0|^2}{\Delta_0}\cos(2kx)+\frac{|V_1|^2}{\Delta_1}\right.$$
$$\left.\times\cos(2(k+\delta k)x+2\phi)\right]+\bar{n}(x)\frac{2\nu_-}{\nu_+}\frac{|V_0|^2}{\Delta_0^2}$$
$$\times\frac{\Gamma_c}{4Dk^2+\gamma}\cos(2kx). \tag{44}$$

A remarkable circumstance is embodied in the fact that for

$$\frac{\nu_-}{\nu_+}\frac{|V_0|^2\Gamma_c}{\Delta_0\gamma}\gg\frac{\hbar|V_1|^2}{\Delta_1 T},\ \frac{\hbar|V_0|^2}{\Delta_0 T}$$

the microscopic oscillations are governed primarily by the light-induced diffusional injection and ejection effect rather than by the light pressure. The relation just written down is most probable for the case of neutral resonant atoms in typical physical situations. Figure 1 illustrates the character of the distribution of the resonant atoms in a macroscopic layer.

Rectified forces in fields with two-dimensional configuration

Let a bichromatic field (1) be formed by the appropriate superposition of traveling light waves intersecting the xy plane and polarized in the direction of the z axis

$$V(\mathbf{r})=V_0\sum_{j=1}^{3}e^{i\phi_j},\quad V_1(\mathbf{r})=V_1\sum_{j=1}^{3}e^{i\phi_j'}, \tag{45}$$

where

$$\phi_j=k\mathbf{n}_j\mathbf{r}+\hat{\phi},\quad \phi_j'=(k+\delta k)\mathbf{n}_j'\mathbf{r}+\hat{\phi}_j',\quad j=1,2,3,$$

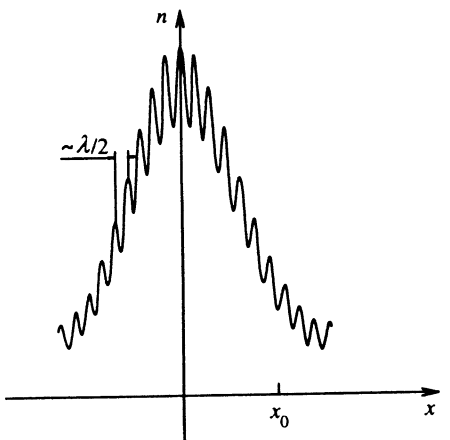

FIG. 1. Spatial dependence of the density of resonant atoms in an individual macroscopic layer.

and $\delta k \ll k$, the angle between the unit wave vectors $\mathbf{n}_j$ is equal to $2\pi/3$, and the system of unit wave vectors $\mathbf{n}'_j$ is "rigidly" rotated about the z axis relative to the vectors $\mathbf{n}_j$ by the small angle $\alpha_0 \ll 1$:

$$\mathbf{n}'_j \simeq \mathbf{n}_j + \alpha_0 [\mathbf{e}_z \mathbf{n}_j], \quad j=1,2,3.$$

A distinguishing feature of such a symmetric spatial configuration of fields is the fact that the radiation flux (equivalently, the momentum density of the field) averaged over the rapid spatial oscillations is zero:[3]

$$\langle \mathbf{J} \rangle = 0. \tag{46}$$

Equations (28) and (35) lead to the following expression for the rectified force in a field of the form (45)

$$\mathbf{F}_R = \frac{8\Gamma_c}{3Dk^2+\gamma} \left| \frac{V_0}{\Delta_0} \right|^2 (\nabla \hat{U} + \text{rot } \mathbf{A}), \tag{47}$$

where

$$\hat{U} = \frac{\hbar |V_1|^2}{\Delta_1} \frac{\alpha}{\alpha_0^2+\alpha^2} u(\mathbf{r}), \quad \mathbf{A} = \frac{\alpha_0}{\alpha} \hat{U} \mathbf{e}_z, \quad \alpha = \frac{\delta k}{k},$$

$$u(\mathbf{r}) = \cos(\bar{k}\mathbf{n}_{12}\mathbf{r} + \psi_{12}) + \cos(\bar{k}\mathbf{n}_{23}\mathbf{r} + \psi_{23}) + \cos(\bar{k}\mathbf{n}_{31}\mathbf{r} + \psi_{31}),$$

$$\mathbf{n}_{jl} = \frac{1}{\sqrt{3(\alpha_0^2+\alpha^2)}} \{[\mathbf{e}_z, (\mathbf{n}_j - \mathbf{n}_l)]\alpha_0 + (\mathbf{n}_j - \mathbf{n}_l)\alpha\},$$

$$\mathbf{n}_{jl}^2 = 1, \quad \psi_{jl} = (\hat{\phi}'_j - \hat{\phi}'_l) - (\hat{\phi}_j - \hat{\phi}_l),$$

$$\bar{k} = \sqrt{3(\alpha_0^2+\alpha^2)} k.$$

Thus, the rectified force has a periodic potential-vortical structure (with spatial period $\sim 1/\bar{k}$), and the ratio of the vortical to the potential component is regulated by the parameter α_0/α. For $\alpha \gg \alpha_0$ (rotation of the vectors $\mathbf{n}'_j$ relative to the vectors $\mathbf{n}_j$ is absent or small) the force field is mainly potential ($|\nabla \hat{U}| \gg |\text{rot } \mathbf{A}|$), and for $\alpha \ll \alpha_0$ it is mainly vortical. It may be readily observed that the spatial symmetry of the force field is analogous to that of its generating electromagnetic field since the angle between any two of the unit vectors $\mathbf{n}_{12}$, $\mathbf{n}_{23}$, and $\mathbf{n}_{31}$ is equal to $2\pi/3$.[1)]

Let us also consider the components of the effective rectified force in the case $\nu_- \neq 0$. From Eqs. (28), (36), and (37) we at once have

$$\mathbf{F}_{1R} = 0, \quad \mathbf{F}_{2R}^{(2)} = 0. \tag{48}$$

The first of these relations is a consequence of the dependence $\hat{N}_1 = \hat{N}_1(|V(\mathbf{r})|^2)$, (which is valid for a field of the type (45) for any relation between γ and Dk^2) and the second is a consequence of Eq. (46). This demonstrates the possibility of eliminating the light-induced drift effect by the appropriate choice of the symmetry of the field configuration. However, to the fourth order in the field amplitude this effect gives a nonvanishing interference contribution to the rectified flux of resonant atoms since $\mathbf{F}_{2R}^{(4)} \neq 0$:

$$\mathbf{F}_{2R}^{(4)} = \frac{24}{\sqrt{3}} \frac{T\Gamma_c \nu_- |V_1|^2 |V_0|^2}{\Delta_0^4 \Delta_1 \nu_+} (\nabla U_c + \text{rot } \mathbf{A}_c), \tag{49}$$

where

$$U_c = \frac{\alpha_0}{\alpha_0^2+\alpha^2} u_1(\mathbf{r}), \quad \mathbf{A}_c = \frac{\alpha}{\alpha_0^2+\alpha^2} u_1(\mathbf{r}) \mathbf{e}_z,$$

$$u_1(\mathbf{r}) = \sin(\bar{k}\mathbf{n}_{12}\mathbf{r} + \psi_{12}) + \sin(\bar{k}\mathbf{n}_{23}\mathbf{r} + \psi_{23}) + \sin(\bar{k}\mathbf{n}_{31}\mathbf{r} + \psi_{31}).$$

Thus we see by comparing expressions (49) and (47) that in the interesting case of large detunings [see Eq. (8)], the force $\mathbf{F}_{1R}^{(4)}$ is only a small correction to the rectified gradient force ($|F_{1R}^{(4)}/F_R| \ll 1$) if the inequality

$$\hbar \Delta_0 \gg T \frac{\nu_-}{\nu_+} \frac{\gamma}{\Delta_0}. \tag{50}$$

is satisfied.

Two-dimensional stratification and rotation of the resonant admixture

To start with, let $\alpha_0 = 0$ and inequality (50) be satisfied. Then from the solution of the averaged equations of transport (32) and (40) we find that in a field of the form (45) the resonant atoms are grouped in the xy plane near the minima of the rectified force potential according to the Boltzmann distribution (43), in which, taking Eq. (47) into account, it is necessary to set

$$U_R(\mathbf{r}) = -U_0 u(\mathbf{r}), \quad U_0 = \frac{8\Gamma_c}{3Dk^2+\gamma} \left| \frac{V_0}{\Delta_0} \right|^2 \frac{\hbar |V_1|^2}{\Delta_1 \alpha}.$$

For $U_0 \gtrsim T$ there arises in the gas a two-dimensional periodic lattice (with period $4\pi/k\alpha = L$) of well-separated swarms of atoms, having the form of cylinders with their axes parallel to the z axis and with characteristic radius x_0:

$$L \gg x_0 \sim \frac{1}{k\sqrt{\alpha}} \left(\frac{\hbar |V_1|^2}{T\Delta_1} \right)^{-1/2} \gg \frac{1}{k}. \tag{51}$$

A qualitatively new situation arises for $\alpha_0 \neq 0$ since there appears a nonzero rectified vortical flux of resonant atoms

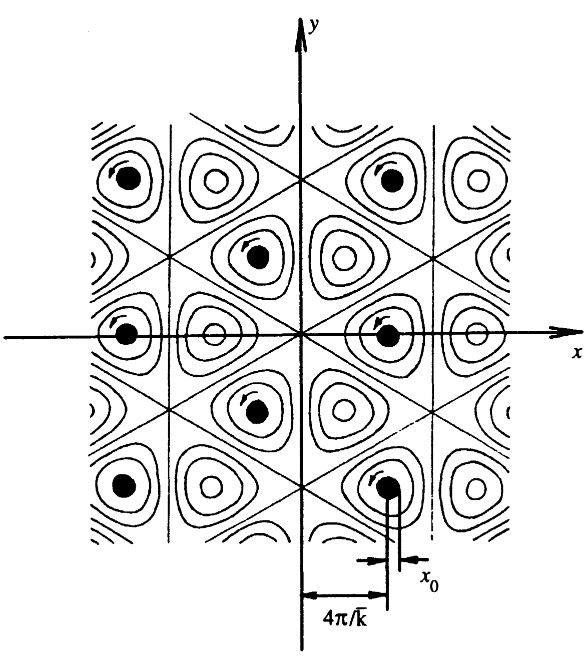

FIG. 2. Typical picture of the periodic stratification of a gas in the xy plane into rotating swarms for a field of the type (45). The thin solid lines are isocontours of the rectified force potential U_R; the arrows indicate the direction of rotation of the swarms. The regions of maximum density are filled in.

$$\mathbf{J}_R=\frac{\text{rot}\,\mathbf{A}_R(\mathbf{r})}{m\nu_+}\,\bar{n}(\mathbf{r}),\quad \mathbf{A}_R=\frac{8\Gamma_c}{3Dk^2+\gamma}\left|\frac{V_0}{\Delta_0}\right|^2\mathbf{A},\qquad(52)$$

where $\mathbf{A}$ is given by formula (47) and $\bar{n}(\mathbf{r})$ is the Boltzmann distribution (43) with potential

$$U_R(\mathbf{r})=-U_0u(\mathbf{r}),$$

$$U_0=\frac{8\Gamma_c}{3Dk^2+\gamma}\left|\frac{V_0}{\Delta_0}\right|^2\frac{\hbar|V_1|^2}{\Delta_1}\frac{\alpha}{\alpha_0^2+\alpha^2}.$$

In the derivation of this solution we made use of the fact that the isocontours of the rectified force potential (47) are simultaneously the vector lines of its vortical component: $\text{rot}\,\mathbf{A}_R\cdot\nabla U_R=0$. Thus, for $U_0\gtrsim T$ there takes place a two-dimensional stratification of the gas into rotating (!) cylinders (see Fig. 2). The characteristic angular velocity of their rotation Ω is given for $\alpha\sim\alpha_0$ by

$$\Omega\sim\omega_R\frac{2|V_1|^2}{\Delta_1\nu_+}\alpha_0.\qquad(53)$$

Let us derive estimates of the radiation intensity and the gas parameters required for light-induced stratification of a gas into a rotating structure. Let $k=5\cdot10^4$ cm^{-1}, $m=100$ a.e.u., $m_\beta=4$ a.e.u., $T=100$ K, $ks\sim6\cdot10^8$ s^{-1}, $\gamma=5\cdot10^7$ s^{-1}, and $\sigma\sim10^{-15}$ cm^2. Then the choice of buffer gas pressure $p\sim5\cdot10^3$ Torr ensures fulfillment of the main conditions (3) and (4) ($\nu\sim3\cdot10^9$ s^{-1}, $k\lambda_c\sim0.2$, $Dk^2\sim\gamma$) and if we set the detunings $\Delta_1\sim2\cdot10^{11}$ Hz and $\Delta_0\sim5\cdot10^{10}$ Hz, and the field intensities $I_0\sim100$ W/cm^{-2} and $I_1>5$ kW/cm^{-2}, then for $\Gamma_c\sim10^{10}$ Hz we have: $\delta k/k\sim10^{-4}$, $U_0>100$ K, $x_0\sim0.1$ cm $\Omega\sim10^4\alpha_0$ s^{-1}, and the period of the macroscopic spatial lattice $L\simeq0.5$ cm. Condition (50) for the prevalence of effects associated with the rectified force is satisfied with sufficient latitude since $\hbar\Delta_0\sim4\cdot10^{-5}$ eV and $\gamma/\Delta_0\sim10^{-3}$. Lowering the temperature of the gas to 10 K, and decreasing the contrast parameter of the atomic lattice $(\bar{n}_{\max}-\bar{n}_{\min})/\bar{n}_{\min}$ from 10 to 0.5 leads to a decrease of the required intensity of the HF field to quite moderate values $I_1\sim200$ W/cm^2. Note that the light-induced macroscopic rotation of the resonant admixture can even cause rotation of the buffer gas if the time of action of the radiation is long enough: $t\sim\tau=\nu^{-1}n_\beta m_\beta/\bar{n}m$. We have not considered this effect here, rather setting $t\ll\tau$ for $n_\beta\gg\bar{n}$.

If condition (50) is fulfilled, as we have seen, large-scale transport of atoms (corresponding to a macroscopic spatial scale $L\gg\lambda$) is governed by the effects of induced light pressure. At the same time, the microstructure of the density and flux distributions of the resonant atoms can be associated for $\nu_-\neq0$ primarily with the light-induced diffusional injection and ejection and light-induced drift effects.

Indeed, for

$$\nu_-/\nu_+\gg\hbar|V_\alpha|^2/\Delta_\alpha T,\qquad \alpha=0,1$$

we obtain from Eq. (39) for a field of the form (45) the following expressions for the rapidly oscillating components of the density and flux of the resonant atoms:

$$\frac{\tilde{n}^{(1)}}{\bar{n}}\simeq\frac{8\nu_-}{\nu_+}\frac{\Gamma_c}{3Dk^2+\gamma}\left|\frac{V_0}{\Delta_0}\right|^2\sum_{j>l}\cos[\phi_j(\mathbf{r})-\phi_l(\mathbf{r})],\qquad(54)$$

$$\tilde{\mathbf{J}}_a=\frac{8T\Gamma_c\nu_-|V_0|^2k}{\nu_+^2\Delta_0^3m}\sum_{j>l}(\mathbf{n}_j+\mathbf{n}_e)\cos[\phi_j(\mathbf{r})-\phi_l(\mathbf{r})].\qquad(55)$$

If the gas is hot enough, the difference of the collision frequencies $\nu_-\neq0$ and is not too small, and the detuning Δ_0 is not too large, then condition (50) is violated and a stratification and rotation of the gas analogous to the one described can occur at the "micro" level as well as at the "macro level," due to effects of the type "light-induced drift" and "light-induced diffusional injection and ejection," i.e., due to the forces $\mathbf{F}_{2R}^{(4)}$ and $\tilde{\mathbf{F}}_1$.[2)]

Note in connection with this that the problem of interference phenomena in the light-induced drift effect analogous to those associated with light pressure[2,3] is, in our opinion, very interesting in its own right and deserve separate study.

6. CONCLUSION

Thus, collisional relaxation is capable of suppressing the damaging effect of velocity selectivity of the rectified gradient force.[3] This makes it possible, with the help of interfering bichromatic fields, a bulk force of significant magnitude (of the order of the induced-light-pressure force) varying only slightly over one wavelength and acting on a resonant admixture contained in a buffer gas.

Collisions with phase jerk lead to the appearance of a rectified radiation force in the fourth order of the expansion over powers of the field amplitude (but not in the sixth, as in the case of an atom unperturbed by collisions[1]).

A bulk, rectified, induced-light-pressure force is capable of acting very efficiently on transport phenomena of a resonant admixture in a buffer gas. Even in the case of different transport collision frequencies of the excited and unexcited atoms in the situation of strong fields and large detunings it can continue to dominate over effects of the type "light-induced diffusional injection and ejection" and "light-induced drift" with regard to the transport of particles over macroscopic spatial scales. It is well known[15,16] that the opposite situation is commonly realized in the case of spontaneous light pressure in a monochromatic field. The spatial structure of the rectified force depends substantially on the configuration of the interfering fields. In the field of a bichromatic one-dimensional standing wave, the rectified force induces rotating, spatially periodic structures in the gas. The magnitude, rotational velocity, and configuration of these structures can be controlled by varying the geometric characteristics (e.g., the angle α_0 in the case of fields of the type (45) and the phase of the acting fields. These effects can be used for spatial localization and separation of small admixtures contained in a gas or plasma, and the creation of light-controlled regular spatial lattices of resonant atoms.

[1)]The system of vectors $\{\mathbf{n}_{jl}\}$ is rotated relative to the vectors $\{\mathbf{n}_j\}$ by the angle $\delta = \arcsin(\alpha_0/\sqrt{\alpha^2+\alpha_0^2})$.

[2)]In such a situation solutions of the averaged equations of the form (52) are defined by the potentials given in Eqs. (49).

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