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Ion drift in a magnetic field under the combined action of LID and light pressure

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Abstract. The effect of magnetic field on the ion drift in a weakly ionised gas under the combined action of light-induced drift (LID) and light pressure is theoretically investigated. It is shown that the imposition of an external magnetic field may give rise to a velocity component of light-induced ion drift orthogonal to the direction of radiation propagation. The effect of light pressure in sufficiently strong magnetic fields is found to prevail over the LID effect, while the reverse is true for weak magnetic fields. The dependence of the ion drift velocity on the frequency detuning drastically changes in the magnetic field when ions experience the Lorenz force. It is predicted that the projection of the ion drift velocity on the direction of radiation propagation should change its sign with increasing magnetic field, and an anomalous LID can be observed.

Keywords *light-induced drift, light pressure, ion drift in a plasma.*

1. Introduction

In the physics of selective radiation action on the translation motion of particles in a buffer medium, the effect of light-induced drift (LID) discovered in 1979 [1, 2] attracts considerable recent attention. The essence of the effect is that particles interacting with a travelling light wave and experiencing collisions with the particles of a buffer gas acquire directional motion.

The LID effect under its optimal conditions exceeds the well-known effect of light pressure by several orders of magnitude [3]. It was theoretically shown that the drift velocity caused by LID under laser excitation can achieve the thermal velocity [4]. It was experimentally demonstrated that LID can force atoms to drift with a velocity of the order of several tens of metres per second [5].

It is easily understood from simple physical considerations that an external magnetic field can exert a strong effect on the LID of atoms, molecules, and ions. Two aspects of this effect can be recognised: the spectral aspect and that of force. The spectral aspect of the magnetic field influence on the LID is due to the Zeeman splitting of the absorption line and manifests itself for any (neutral or charged) particles. The effect of force on the LID induced by the magnetic field manifests itself only for ions and is caused by the Lorenz force acting on the ions drifting in the magnetic field.

Up to the present time, the effect of the magnetic field on the LID has been little studied and is discussed only in Refs [6 – 8]. The spectral aspect of the effect of the magnetic field on the LID was experimentally demonstrated in Ref. [6], which was concerned with the spatial localisation and accumulation of neutral atoms. The force aspect of the magnetic field effect on the LID of ions was theoretically investigated in Refs [7, 8]. Unfortunately, the results of neutral atoms are much higher than their intercollision frequencies. Let the radiation in the form of a travelling monochromatic wave be resonantly absorbed on the $m - n$ transition between the ground (n) and first excited (m) ion levels. Below, we will study only the force action of the magnetic field on the ion drift and therefore restrict ourselves to the consideration of the simplest case, when the Zeeman splitting of the absorption line can be neglected. Such is the case, for instance, for a simple

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* In Refs [7, 8], the influence of the magnetic field on the ion drift velocity $\mathbf{u}(\mathbf{r}, t)$ (in the notation of Refs [7, 8]), which enters in the ion flux equation (Eqn (8) in Ref. [7] and Eqn (1) in Ref. [8]), was not taken into account. The ion flux equation analysed in Refs [7, 8] is actually a modification of the first equation of the system (7) of the present work (for $a = 0$ and $\xi = 0$) and is of the following form in the notation of the present work: $\nu_n \mathbf{j} + (\nu_m - \nu_n)\mathbf{j}_0 = -((\bar{v}^2/2)\nabla N + \omega_c[\mathbf{j}\mathbf{h}]$ where the flux $\mathbf{j}_0$ is independent of the magnetic field, unlike the flux $\mathbf{j}_m$ of Ref. [7].

Zeeman effect when the radiation polarised linearly along the magnetic field B propagates transversely to B.

Under these conditions, the interaction of radiation with ions, taking recoil effect into account is described by the equations for the density matrix [9]:

$$\left[\frac{d}{dt} + \Gamma_m\right]\rho_m(\boldsymbol{v}) = S_m(\boldsymbol{v}) + NP(\boldsymbol{v} - \xi),$$

$$\frac{d}{dt}\rho_n(\boldsymbol{v}) = S_n(\boldsymbol{v}) + \hat{\Gamma}_m\rho_m(\boldsymbol{v}) - NP(\boldsymbol{v} + \xi), \qquad (1)$$

$$\left[\frac{d}{dt} + \frac{\Gamma_m}{2} - i(\Omega - k\boldsymbol{v})\right]\rho_{mn}(\boldsymbol{v})$$

$$= S_{mn}(\boldsymbol{v}) + iG[\rho_n(\boldsymbol{v} - \xi) - \rho_m(\boldsymbol{v} + \xi)],$$

where

$$\frac{d}{dt} \equiv \frac{\partial}{\partial t} + \boldsymbol{v}\frac{\partial}{\partial r} + a_i\frac{\partial}{\partial \boldsymbol{v}}; \quad a_i = \frac{eE}{M} + \omega_c[\boldsymbol{v}h]; \quad \omega_c = \frac{eB}{Mc};$$

$$\hat{\Gamma}_m\rho_m(\boldsymbol{v}) = \frac{\Gamma_m}{4\pi}\int\rho_m(\boldsymbol{v} + 2\xi n_r)dn_r; \qquad (2)$$

$$NP(\boldsymbol{v}) = -2\mathrm{Re}[G\rho_{mn}(\boldsymbol{v})],$$

$$\xi = \frac{\hbar k}{2M}; \quad |G|^2 = \frac{B_{nm}I}{2\pi}; \quad \Omega = \omega - \omega_{mn};$$

$h = B/B$; $\rho_i(\boldsymbol{v})$ is the velocity distribution function for the ions at a level $i = m, n$; N is the total ion density; $S_m(\boldsymbol{v})$, $S_n(\boldsymbol{v})$, and $S_{mn}(\boldsymbol{v})$ are the ion collision integrals; ω and k are the radiation frequency and wave vector, respectively; Γ_m is the spontaneous decay probability of the excited m state; $\hat{\Gamma}_m\rho_m(\boldsymbol{v})$ is the integral operator describing the radiative transition of particles from the excited level m to the ground level n, taking into account that the particle velocity changes due to the recoil effect upon spontaneous emission; n_r is the unit vector defining the direction of spontaneous emission; ω_{mn} is the $m - n$ transition frequency; B_{nm} is the Einstein coefficient for the $m - n$ transition; I is the emission intensity; ω_c is the cyclotron frequency of ion gyration; e is the elementary charge; M is the ion mass; B is the magnetic field induction; and E is the internal electric field intensity inside the medium.

The electric field E in the medium may arise from the directional motion of the ions as a whole due to LID and light pressure. In this case, two different situations are possible. When the concentration of charged particles is not high enough for the ionised gas to exhibit plasma properties ($r_d \gg L$, where r_d is the Debye radius and L is the characteristic size of the system), the electrons exert no effect on the ion drift and the field E in Eqns (1) can be neglected.

When the concentration of charged particles is high enough for the ionised gas to exhibit plasma properties ($r_d \ll L$), the directional ion motion should bring about (owing to the quasineutrality condition for plasma) a directional electron motion. This gives rise to the electric field, which compensates for the friction of electrons against buffer particles.

Therefore, the electron motion in a plasma is related to the ion motion under the action of the electric field E, and Eqns (1) should be supplemented with the equation for the electron distribution function $\rho_e(\boldsymbol{v})$:

$$\left[\frac{\partial}{\partial t} + \boldsymbol{v}\frac{\partial}{\partial r} + a_e\frac{\partial}{\partial \boldsymbol{v}}\right]\rho_e(\boldsymbol{v}) = S_e(\boldsymbol{v}),$$

$$a_e = -\frac{eE}{m_e} - \omega_e[\boldsymbol{v}h], \qquad (3)$$

where $\omega_e = eB/m_ec$ is the electron cyclotron frequency; m_e is the electron mass; and $S_e(\boldsymbol{v})$ is the electron collision integral.

Under the condition when the velocity distribution functions of particles are close to the Maxwell function, kinetic equations (1) and (3) can be solved by the Grad method (the method of moments) [10]. We will solve the problem employing the simplest approximation of the Grad method. This approximation implies that the velocity dependence of the elements of the density matrix in Eqns (1) and (3) can be represented as the sum of the Maxwell distribution $W(\boldsymbol{v})$ and an antisymmetric correction:

$$\rho_i(\boldsymbol{v}) = \left[N_i + \frac{2}{\bar{v}_\alpha^2}\boldsymbol{v}j_i\right]W(\boldsymbol{v}), \quad i = m, n, mn, e, \qquad (4)$$

where $N_i = \int\rho_i(\boldsymbol{v})d\boldsymbol{v}$; $j_i = \int\boldsymbol{v}\rho_i(\boldsymbol{v})d\boldsymbol{v}$; N_m, N_n and j_m, j_n are the densities and the ion fluxes, respectively, in the states m and n; N_e and j_e are the electron density and flux, respectively; $\bar{v}_\alpha$ is the most probable velocity of particles of the α type [$\bar{v}_\alpha = \bar{v} = (2k_BT/M)^{1/2}$ for ions ($i = m, n, mn$) and $\bar{v}_\alpha = \bar{v}_e = (2k_BT/m_e)^{1/2}$ for electrons ($i = e$), where T is the temperature and k_B is the Boltzmann constant]. Expression (4) is applicable when the radiation–ion interaction is weakly sensitive to velocities, i.e., for a homogeneous broadening of the absorption line $\Gamma > k\bar{v}$, where Γ is the homogeneous half-width of the absorption line of ions and $k\bar{v}$ is the Doppler width*.

For the nondiagonal collision integral in Eqns (1), we will use the approximation $S_{mn}(\boldsymbol{v}) = -(\Gamma - \Gamma_m/2)\rho_{mn}(\boldsymbol{v})$ commonly employed in nonlinear spectroscopy, which implies that collisions cause a complete dephasing of the oscillating dipole moment [9].

Inelastic collisions are insignificant for the problem involved (the effective ionisation and recombination rates are low compared to the elastic collision rates), and we will take into account only the elastic collisions of ions and electrons with buffer particles (neutral atoms). In elastic collisions, the relation $\int S_i(\boldsymbol{v})d\boldsymbol{v} = 0$ ($i = m, n, e$) holds for the nondiagonal collision integrals in Eqns (1) and (3), which means the conservation of the total number of particles (the ions in the states m or n and the electrons).

The first moment of the diagonal collision integrals in the approximation (4) obeys the relation $\int\boldsymbol{v}S_i(\boldsymbol{v})d\boldsymbol{v} = -\nu_i j_i$ ($i = m, n, e$), where ν_i is the average transport collision frequency [2]. For ions ($i = m, n$), the average transport frequency and the diffusion coefficient D_i for ions in the state i are related by a simple formula: $D_i = \bar{v}^2/2\nu_i$. For electrons ($i = e$), the diffusion coefficient is $D_e = \bar{v}_e^2/2\nu_e$.

*Approximation (4) is valid for an arbitrary ratio between Γ and $k\bar{v}$ in the case of broadband radiation with a spectral shape smooth enough within the absorption line width.

Integration of Eqns (1) with respect to $\boldsymbol{v}$ leads, in view of expressions (4), to the equations:

$$\frac{\partial N}{\partial t} + \operatorname{div}\mathbf{J} = 0,$$

$$\left(\frac{\partial}{\partial t} + \Gamma_m\right)N_m + \operatorname{div}\mathbf{j}_m = NP, \tag{5}$$

$$\left(\frac{\partial}{\partial t} + \Gamma - i\Omega\right)N_{mn} + i\mathbf{k}\,\mathbf{j}_{mn} + \operatorname{div}\mathbf{j}_{mn} = iG(N - 2N_m),$$

where

$$P \equiv \int P(\boldsymbol{v})\,d\boldsymbol{v} = -\frac{2}{N}\,\mathrm{Re}(iG^{*}N_{mn}); \qquad N = N_m + N_n; \quad \mathbf{J} = \mathbf{j}_m + \mathbf{j}_n; \tag{6}$$

P is the number of radiation absorption events per unit time by an ion; and $\mathbf{J}$ is the total ion flux.

By multiplying Eqns (1) by $\boldsymbol{v}$ and then integrating them with respect to $\boldsymbol{v}$, we obtain

$$\left(\frac{\partial}{\partial t} + \nu_n\right)\mathbf{J} + (\nu_m - \nu_n)\mathbf{j}_m = \left(\boldsymbol{\alpha} - \frac{\bar{v}^2}{2}\nabla\right)N + \omega_c[\mathbf{J}\mathbf{h}] + 2NP\boldsymbol{\xi},$$

$$\left(\frac{\partial}{\partial t} + \Gamma_m + \nu_m\right)\mathbf{j}_m = \left(\boldsymbol{\alpha} - \frac{\bar{v}^2}{2}\nabla\right)N_m + \omega_c[\mathbf{j}_m\mathbf{h}] - 2\mathrm{Re}(iG^{*}\mathbf{j}_{mn}) + NP\boldsymbol{\xi}, \tag{7}$$

$$\left(\frac{\partial}{\partial t} + \Gamma - i\Omega\right)\mathbf{j}_{mn} = \left[\boldsymbol{\alpha} - \frac{\bar{v}^2}{2}(i\mathbf{k} + \nabla)\right]N_{mn} + \omega_c[\mathbf{j}_{mn}\mathbf{h}] + iG(\mathbf{J} - 2\mathbf{j}_m) + iGN\boldsymbol{\xi},$$

where $\boldsymbol{\alpha} = e\mathbf{E}/M$; Ω is the radiation frequency detuning.

Similarly, for the zero and first momenta of kinetic equation (3) we obtain

$$\frac{\partial N_e}{\partial t} + \operatorname{div}\mathbf{j}_e = 0,$$

$$\left(\frac{\partial}{\partial t} + \nu_e\right)\beta\mathbf{j}_e + \left(\boldsymbol{\alpha} + \frac{\beta\bar{v}_e^2}{2}\nabla\right)N_e + \omega_c[\mathbf{j}_e\mathbf{h}] = 0, \tag{8}$$

where $\beta = m_e/M$ is the electron-to-ion mass ratio.

By definition, the ion drift velocity is $\mathbf{u} \equiv \mathbf{J}/N$ and is found from the system of equations (5), (7), and (8).

3. Ion drift under gas conditions

Consider first the case when the concentration of charged particles is not high enough for the ionised gas to acquire plasma properties ($r_d \gg L$). In this case, electrons have no effect on the ion drift, and one can put $\boldsymbol{\alpha} = 0$ in Eqns (7). Then, under stationary and spatially uniform conditions for radiation propagating perpendicular to the magnetic field (for $\mathbf{k} \perp \mathbf{B}$), we find from the solution of the system of equations (5) and (7) that the ion drift velocity is equal to the sum of two mutually perpendicular components $\mathbf{u}_k$ and $\mathbf{u}_\perp$:

$$\mathbf{u} = \mathbf{u}_k + \mathbf{u}_\perp. \tag{9}$$

Here, the component $\mathbf{u}_k$ is parallel to the wave vector $\mathbf{k}$ and the component $\mathbf{u}_\perp$ is perpendicular to $\mathbf{k}$ and $\mathbf{B}$:

$$\mathbf{u}_k = \frac{\mathbf{k}}{k}u_k, \quad \mathbf{u}_\perp = \mathbf{n}u_\perp, \quad \mathbf{n} = \frac{[\mathbf{k}\mathbf{B}]}{kB}. \tag{10}$$

The drift velocity can also be represented as the sum of the drift velocities $\mathbf{u}_L$ and $\mathbf{u}_r$ caused by the effects of LID ($\mathbf{u}_L$) and light pressure ($\mathbf{u}_r$):

$$\mathbf{u} = \mathbf{u}_L + \mathbf{u}_r. \tag{11}$$

For an arbitrary radiation intensity, the formulas for the projections u_{Lk}, u_{rk}, $u_{L\perp}$, and $u_{r\perp}$ of the drift velocities $\mathbf{u}_L$ and $\mathbf{u}_r$ on the directions of $\mathbf{k}$ and $\mathbf{n}$ are cumbersome. For a low radiation intensity, i.e. when $I \ll \Gamma\Gamma_m/B_{mn}$, these formulas are considerably simpler. In this case, these formulas, correct to small terms of the order of $(k\bar{v}/\Gamma)^2 \ll 1$, $k\xi/\Gamma \ll 1$, take the form:

$$u_{Lk} \equiv \frac{\mathbf{k}}{k}\mathbf{u}_L = u_0\left\{1 - \frac{\omega_c^2}{\nu_n(\Gamma_m + \nu_m)}\left[1 + \frac{(3\Gamma^2 + \omega_c^2 - \Omega^2)(\Gamma_m + \nu_m + \nu_n)}{2\Gamma(\Gamma^2 + \Omega^2)}\right]\right\}, \tag{12}$$

$$u_{L\perp} \equiv \mathbf{n}\mathbf{u}_L = u_0\,\frac{\omega_c(\Gamma_m + \nu_m + \tilde{\nu}_n)}{\nu_n(\Gamma_m + \nu_m)}\left\{1 + \frac{(3\Gamma^2 + \omega_c^2 - \Omega^2)\left[(\Gamma_m + \nu_m)\tilde{\nu}_n - \omega_c^2\right]}{2\Gamma(\Gamma^2 + \Omega^2)(\Gamma_m + \nu_m + \tilde{\nu}_n)}\right\}, \tag{13}$$

$$u_{rk} \equiv \frac{\mathbf{k}}{k}\mathbf{u}_r = u_{0r}\left\{1 + (\nu_n - \nu_m)\left[2\left(\Gamma_m + \nu_m + \frac{\omega_c^2}{\Gamma_m + \nu_m}\right)\right]^{-1}\left[f_1 - \frac{\omega_c^2 f_2}{\nu_n(\Gamma_m + \nu_m)}\right]\right\}, \tag{14}$$

$$u_{r\perp} \equiv \mathbf{n}\mathbf{u}_r = u_{0r}\,\frac{\omega_c}{\nu_n}\left\{1 + (\nu_n - \nu_m)\left[2\left(\Gamma_m + \nu_m + \frac{\omega_c^2}{\Gamma_m + \nu_m}\right)\right]^{-1}\left[f_1 + \frac{\tilde{\nu}_n f_2}{\Gamma_m + \nu_m}\right]\right\}, \tag{15}$$

where

$$u_0 = \bar{v}\tau_\sigma P\varphi; \quad u_{0r} = \frac{2\xi P}{\tilde{\nu}_n + \omega_c^2/\nu_n}; \quad P = \frac{2|G|^2\Gamma}{\Gamma^2 + \Omega^2};$$

$$\tau_\sigma = (\nu_n - \nu_m)\left(\tilde{\nu}_n + \frac{\omega_c^2}{\nu_n}\right)^{-1}\left(\Gamma_m + \nu_m + \frac{\omega_c^2}{\Gamma_m + \nu_m}\right)^{-1};$$

$$\varphi = \frac{\Omega k\bar{v}(\Gamma^2 + \Omega^2)}{q}; \quad f_1 = f - \frac{\omega_c^2 g}{(\Gamma_m + \nu_m)^2}; \quad f_2 = f + g; \tag{16}$$

$$f = 1 + \frac{(\Gamma^2 + \Omega^2)(\Gamma^2 + \omega_c^2 + \Omega^2)}{q};$$

$$g = \frac{(\Gamma^2 + \Omega^2)(\Gamma^2 + \omega_c^2 - \Omega^2)(\Gamma_m + \nu_m)}{q\Gamma};$$

$$q = \left[\Gamma^2 + (\Omega + \omega_c)^2\right]\left[\Gamma^2 + (\Omega - \omega_c)^2\right].$$

In the above formulas, $\tilde{\nu}_n = \nu_n$. The notation $\tilde{\nu}_n$ was introduced here to unify formulas (12)–(16), which will be also valid under plasma conditions if $\tilde{\nu}_n$ from (20) is substituted in them, as follows from Section 4.

The component u_L of the drift velocity u is caused by the LID effect. This component is independent of the light pressure (independent of the ion recoil velocity 2ξ upon absorption of a photon) and $u_L \neq 0$ only when the ion transport collision frequencies for the initial and final transition levels are different (for $\nu_m \neq \nu_n$), as should be the case with LID. In the absence of a magnetic field (for $\omega_c = 0$), from (12) and (13) there follows a well-known formula for the LID velocity for a homogeneous broadened absorption line [1, 2]. The component u_L of the drift velocity u is caused by light pressure: u_r is proportional to the ion recoil velocity 2ξ upon absorption of a photon.

The drift velocities calculated from expressions (12)–(15) are plotted in Figs 1–3. Throughout this paper, as the unit of measurement of velocity we adopted the quantity $u_R = 2\xi P_0/\nu_n$. This quantity is equal to the maximum (for $\Omega = 0$) ion drift velocity u_r induced by the light pressure in the absence of a magnetic field and for $\nu_m = \nu_n$ (here, P_0 is the value of P for $\Omega = 0$). The ratio between the maximum velocities u_L and u_r in the absence of the magnetic field is characterised by the parameter A:

$$\frac{|(u_L)_{\omega_c \to 0}|_{max}}{|(u_r)_{\omega_c \to 0}|_{max}} = \frac{1}{6}\frac{k\bar{v}^2}{\xi\Gamma}\frac{|\nu_n - \nu_m|}{\Gamma_m + \nu_n} = A. \quad (17)$$

For ions with a mass $M \approx 20$ amu in the optical spectral region for a temperature $T \approx 1000$ K and $k\bar{v}/\Gamma \approx 0.2$, the parameter $k\bar{v}^2/\xi\Gamma \approx 10^4$. We used this value in our numerical calculations. For the values of parameters employed in the calculation of the curves plotted in Figs 1 and 2, $A \approx 150$.

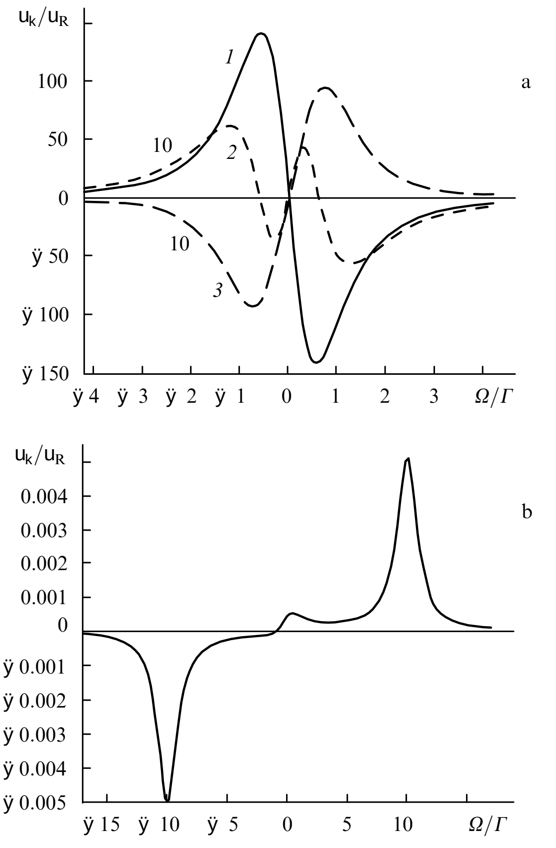

Figure 1. Projection u_k of the drift velocity u on the radiation direction as a function of the radiation frequency detuning Ω for different values of the magnetic field $\omega_c/\nu_n = 5 \times 10^{-3}$ (*1*), 0.9 (*2*), and 3.75 (*3*) (a) and $\omega_c/\Gamma = 10$ (b); $k\bar{v}^2/\xi\Gamma = 10^4$, $\nu_n/\Gamma = 0.2$, $(\nu_m - \nu_n)/\nu_n = 0.1$, and $\Gamma_m/\nu_n = 5 \times 10^{-2}$.

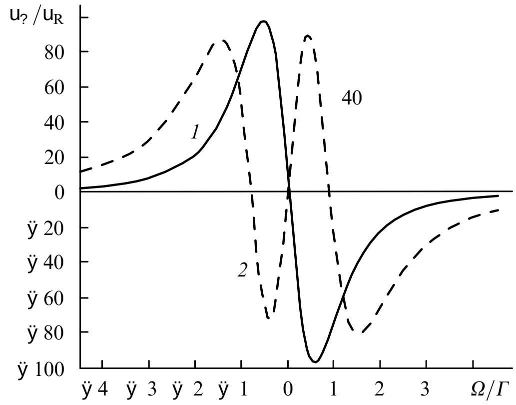

Figure 2. Projection $u_{\perp}$ of the drift velocity u on the direction n as a function of the radiation frequency detuning for different values of the magnetic field $\omega_c/\nu_n = 0.6$ (*1*) and 3.75 (*2*); $k\bar{v}^2/\xi\Gamma = 10^4$, $\nu_n/\Gamma = \nu_n/\Gamma = 0.2$, $(\nu_m - \nu_n)/\nu_n = 0.1$, and $\Gamma_m/\nu_n = 5 \times 10^{-2}$.

One can see from Fig. 1a that the projection u_k of the drift velocity u on the wave vector k changes its sign with increasing the magnetic field. The sign changes when the cyclotron frequency $\omega_c \approx [\nu_n(\Gamma_m + \nu_m)]^{1/2}$. For parameters taken for the calculation of the curves in Fig. 1a, the velocity u_k is determined by the LID effect: $u_k \approx u_{Lk}$. Curves (*1*) and (*3*) correspond to the conventional LID with a characteristic dispersion-like frequency dependence of the drift velocity $u_{Lk}(\Omega)$, which intersects the Ω/Γ axis once for $\Omega = 0$. Curve (*2*) intersects the above axis three times and corresponds to the so-called anomalous LID [11] with a strong departure of the frequency dependence of the drift velocity $u_{Lk}(\Omega)$ from the dispersion-like curve. The anomalous LID appears when the cyclotron frequency of ion gyration is $\omega_c \approx [\nu_n(\Gamma_m + \nu_m)]^{1/2}$. An analysis shows that the interval $\Delta\omega_c$ of cyclotron frequency values where the anomalous LID is observed is equal to $\approx 0.2\nu_n$.

As the magnetic field is further increased (for $\omega_c \gtrsim 5\Gamma$), the dependence $u_k(\Omega)$ first exhibits two peaks and then three (Fig. 1b) peaks with the FWHM 2Γ. Two antisymmetrically located side peaks are at a distance of $|\Omega| = \omega_c$ from the absorption line centre and are caused by the LID effect. The central peak (for $\Omega = 0$) arises from the drift induced by light pressure. The amplitudes $[|(u_{Lk})_{|\Omega| = \omega_c}| \propto 1/\omega_c^3]$ of the two side peaks and the amplitude $[(u_{rk})_{\Omega = 0} \propto 1/\omega_c^2]$ of the central peak decrease with increasing magnetic field. Therefore, the light pressure 'competes' with LID or even prevails over this effect upon the ion drift along the radiation direction in a sufficiently strong magnetic field.

Fig. 2 shows the dependence of the projection $u_{\perp}$ of the drift velocity on the radiation frequency detuning. For $\omega_c \lesssim \Gamma$, the velocity $u_{\perp}$ is determined by the LID: $u_{\perp} \approx u_{L\perp}$. The absolute (in Ω and ω_c) maximum of velocity $u_{\perp}$ is attained for $\omega_c \approx \nu_n/2$ and is close to the absolute velocity maximum u_k, which is attained for $\omega_c \to 0$. For a cyclotron frequency of ion gyration $\omega_c \gtrsim 3\nu_n$, there occurs an anomalous LID transverse to the wave vector $\mathbf{k}$ (curve *2* in Fig. 2).

For $v_m < v_n$, in strong magnetic fields ($\omega_c \gg \Gamma$) there appears an interesting feature in the dependence $u_{rk}(\Omega)$: in the vicinity of radiation frequency detuning $|\Omega| = \omega_c$, the velocity u_{rk} becomes negative (Fig. 3), i.e., the light pressure causes the ion drift in the opposite direction to the radiation propagation ('negative' light pressure)*. It follows from (14) that for $\omega_c \gg \Gamma$ the ratio between maximum velocities u_{rk} for $\Omega = 0$ and $|\Omega| = \omega_c$ is determined by the relative difference of transport collision frequencies of ions in the ground and excited states with buffer particles

$$\frac{(u_{rk})_{|\Omega| = \omega_c}}{(u_{rk})_{|\Omega| = 0}} = \frac{v_m - v_n}{2(v_m + v_n)}. \tag{18}$$

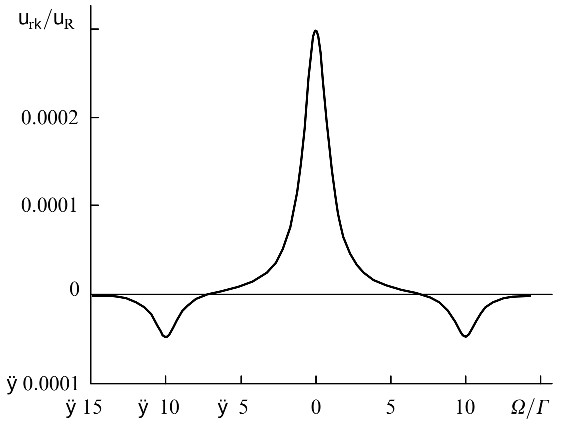

Figure 3. Projection u_{rk} of the drift velocity $\mathbf{u}_r$ caused by light pressure as a function of the radiation frequency detuning for $\omega_c/\Gamma = 10$, $v_n/\Gamma = 0.2$, $(v_n - v_m)/v_n = 0.5$, and $\Gamma_m/v_n = 5 \times 10^{-2}$

4. Ion drift under plasma conditions

Consider now the case when the concentration of charged particles is high enough for the ionised gas to exhibit plasma properties ($r_d \ll L$). In this case, the system of Eqns (5) and (7) should be solved in combination with Eqns (8).

From the continuity equation for ions and electrons taking into account the plasma quasi-neutrality $N_e = N$, the condition $\operatorname{div} \mathbf{J} = \operatorname{div} \mathbf{j}_e$ follows, which relates the ion and electron fluxes flowing into every volume element. Consider a cylindrical plasma column uniform in azimuth φ and along the z axis located in a uniform magnetic field directed along the z axis. Let a cylindrical monochromatic wave with a wave vector $\mathbf{k}$ perpendicular to the z axis diverge radially from a radiation source extended along the z axis. Then, the velocity of light-induced ion drift depends only on the radius ρ. From the condition $\operatorname{div} \mathbf{J} = \operatorname{div} \mathbf{j}_e$, there follows the equality of the radial components of the fluxes $J_\rho = j_{e\rho}$, i.e., the drift of ions and electrons along the $\mathbf{k}$ direction will be ambipolar. Under stationary conditions, the internal electric field $\mathbf{E}$, which arises in the plasma due to the ion drift, is vortex-free, and therefore the azimuthal field component $E_\varphi = 0$. Therefore, in the case under study the condition $\operatorname{div} \mathbf{J} = \operatorname{div} \mathbf{j}_e$ is equivalent to the conditions

$$j_{ek} = J_k, \quad \boldsymbol{\alpha} = \boldsymbol{\alpha}_k, \tag{19}$$

where the symbol k denotes the vector component directed along $\mathbf{k}$.

Individual parts of the cylindrical wave that are small compared to the distance to the radiation source behave as plane waves with a constant radiation intensity. The light-induced drift of particles in these parts can be treated as a drift under spatially uniform conditions.

Under stationary and spatially uniform conditions at a low radiation intensity, we find from the system of Eqns (5), (7), and (8) taking (19) into account, that the ion drift velocity in plasma is given by previous expressions (9)–(16) after substitution of the quantity $\tilde{v}_n$:

$$\tilde{v}_n = v_n + \beta v_e + \frac{\omega_c^2}{\beta v_e}. \tag{20}$$

into them.

In expressions (12) and (14) for u_{Lk} and u_{rk}, the quantity $\tilde{v}_n$ appears only in the factor $1/(\tilde{v}_n + \omega_c^2/v_n)$, and, therefore, the dependences of these velocities on the radiation frequency detuning Ω remain the same as in gas. Only the velocities become lower. This reduction is due to the ambipolar nature of the drift along the radiation direction and the decelerating action of electrons on the ion drift.

The dependence $u_\perp(\Omega)$ changes in passing from gas to plasma. For $\omega_c \lesssim \Gamma$, the dependence $u_\perp(\Omega) \approx u_{L\perp}(\Omega)$ is dispersion-like [the $u_\perp(\Omega)$ line shape is similar to the shape of curve (*1*) in Fig. 2]. As the magnetic field increases (for $\omega_c \gtrsim 5\Gamma$), two antisymmetrically located peaks with the FWHM 2Γ appear in the dependence $u_\perp(\Omega) \approx u_{L\perp}(\Omega)$ at a distance $|\Omega| = \omega_c$ from the absorption line centre.

From Eqns (8) it is easy to find taking (19) into account the ambipolar electric field $\mathbf{E}$, which is automatically produced in the plasma to equalise the oppositely charged particle fluxes along the radiation direction $\mathbf{k}$ (stationary and spatially uniform conditions):

$$E = -\frac{u_k}{\mu_{e\perp}}, \quad \mu_{e\perp} = \mu_e \left(1 + \frac{\omega_e^2}{v_e^2}\right)^{-1}, \quad \mu_e = \frac{e}{m_e v_e}. \tag{21}$$

Here, $\mu_{e\perp}$ is the electron mobility in the direction perpendicular to the magnetic field [10]; and μ_e is the electron mobility in the absence of a magnetic field.

Consider the dependence of E on the magnetic field. When $v_m \neq v_n$, the drift velocity u_k in expression (21) in the limit of weak magnetic fields is determined by the LID, and the maximum intensity of the ambipolar electric field (attained for $|\Omega| = \Gamma/\sqrt{3}$) is

$$|(E)_{\omega_c \to 0}|_{\max} \simeq \frac{A u_R}{\mu_e}. \tag{22}$$

In the limit of strong magnetic fields, the drift velocity u_k is determined by light pressure, and the maximum value of E (attained for $\Omega = 0$) is

$$|(E)_{\omega_c \to \infty}|_{\max} \simeq \frac{u_R}{\mu_e} \frac{v_n}{\beta v_e} = \frac{\hbar k P_0}{e}. \tag{23}$$

*A similar effect for neutral atoms (the light-induced particle drift in opposite direction to the radiation propagation) was predicted in Ref. [12] and was called 'negative' light pressure. For neutral atoms, this effect is possible in an intense light field for $v_m > v_n$ and is manifested in a pure form for a zero radiation frequency detuning $\Omega = 0$.

Therefore, the intensity of ambipolar electric field changes by a factor of $A\beta\nu_e/\nu_n$ in passing from weak magnetic fields to strong ones. If we set for an estimate $\beta\nu_e/\nu_n \sim 10^{-2}$ and $A \sim 150$ [see formula (17) and the comment to it], then $A\beta\nu_e/\nu_n \sim 1$, and hence the intensity of ambipolar electric field is almost independent of the magnetic-field strength for the given parameter values. From expressions (23) for $P_0 \sim 10^7$ s^{-1} we obtain the estimate $|E| \sim 10^{-3}$ V cm^{-1} in the optical spectral range.

5. Conclusions

We have studied the force action of an external magnetic field on the light-induced ion drift under the conditions when this action is maximal and is manifested in a 'pure' form (the Zeeman splitting of absorption line is absent). The force action is maximal when the magnetic field B is perpendicular to the direction of radiation propagation k, and it is in this case that this action can be observed in a 'pure' form (the line splitting is absent in the case of a simple Zeeman effect, when $k \perp B$ and the radiation is linearly polarised along B). The expressions for the ion drift velocity u in the field of monochromatic radiation obtained in this work are valid for a homogeneous broadening of the absorption line ($\Gamma > k\bar{v}$).

Note that expressions (12)–(15) describing the ion drift under the action of a travelling monochromatic light wave can be naturally generalised to the case of broadband radiation with an arbitrary spectral intensity $I(\omega)$. For this purpose, it is sufficient to replace the radiation intensity I, which appears in the factor $|G|^2$ in expression (16) for P, by the spectral intensity $I(\omega)$ dependent on the radiation frequency ω and then integrate infinitely over ω the obtained modified expressions (12)–(15) for the drift velocity $u(\omega)$. The expressions obtained in this way will be also valid for a Doppler ($k\bar{v} \gg \Gamma$) broadening of the absorption line in the case of broadband radiation with a sufficiently smooth spectral shape within the absorption line width. Such conditions (the Doppler broadening and a sufficiently smooth shape of the emission spectrum within the absorption line width) are typical, for instance, of stellar atmospheres.

To observe the reversal of direction of the drift along k and an anomalous LID, magnetic fields are required which provide a cyclotron frequency of ion gyration $\omega_c \gtrsim [\nu_n(\Gamma_m + \nu_m)]^{1/2}$. Hence, it follows that the lower are the gas pressure and the spontaneous decay rate of excited ion state, the lower is the magnetic field required for the experimental observation of these effects. For a transport ion collision frequency $\nu_n \sim 10^5$ s^{-1} (this corresponds to a gas pressure of $\sim$0.01 Torr) and a radiative decay probability $\Gamma_m \sim 10^7$ s^{-1}, we obtain the estimate $\omega_c \sim 10^6$ s^{-1}. For the ion mass $M \sim 20$ amu, this cyclotron frequency is attained in moderate magnetic fields $B \sim 2 \times 10^3$ G.

The component of drift velocity transverse to the direction of radiation propagation arises in arbitrarily weak magnetic fields. It can be estimated by the formula $|u_\perp| \sim (\omega_c/\nu_n)|u_d|$, where u_d is the ion drift velocity in the absence of a magnetic field. The transverse drift velocity can reach the value $|u_d|$ even in relatively weak magnetic fields ($B \sim 200$ G for $\nu_n \sim 10^5$ s^{-1} and ion mass $M \sim 20$ amu).

Under laboratory conditions, the light-induced ion drift can manifest itself as an electric current (a light-induced current [13]). The potential difference $V \sim |E|L$ will be produced across the ends of a cell with a weakly ionised gas, where L is the cell length and E is the ambipolar electric field inside the cell arising due to the light-induced ion drift. For $|E| \sim 10^{-3}$ V cm^{-1} and $L \sim 10$ cm, the potential difference $V \sim 10^{-2}$ V will appear across the cell ends. The electric current $I \sim V/R$ will flow through a conductor connecting the opposite cell ends, where R is the internal plasma resistance. Since $R \sim L/eNS\mu_e$, where S is the cross-sectional cell area, we obtain, taking (21) into account, an estimate $I \sim |u_k|eNS$. This gives, for the drift velocity $|u_k| \sim 1$ cm s^{-1}, the ion density $N \sim 10^{11}$ cm^{-3}, and $S \sim 1$ cm^2, the value of current $I \sim 10^{-8}$ A.

The results obtained in this work may be of interest for astrophysical applications in connection with the phenomenon of chemically peculiar stars, which now is under discussion in the scientific literature [14–17]. One of the main hypotheses explains the origin of chemical composition anomalies of all peculiar stars by the separation of chemical elements in their atmospheres due to the selective drift of atoms and ions caused by stellar radiation [14–17]. Both light pressure [15–17] and LID [17–19] were considered as the cause of drift in stellar atmospheres. Among the chemically peculiar stars are the so-called magnetic stars [14–17] with strong (up to 3×10^4 G) large-scale magnetic fields, which are predominantly dipole in nature. As shown in this work, the magnetic field changes drastically the ion drift pattern and can therefore strongly affect the separation of chemical elements in the atmospheres of magnetic stars.

When considering the ion drift in a magnetic field, we assumed it to be given and neglected the back action of the drift on this field. The drift motion may be accompanied (for a specific geometrical configuration) by the appearance of a closed electric current. This gives rise to a magnetic perturbation $\delta B \sim 4\pi jL/c$, where j is the current density and L is the characteristic length of the process. For laboratory plasmas (in small-sized media), the magnetic perturbation caused by the ion drift is negligible. However, in media with large characteristic dimensions (astrophysical objects), the electric current caused by light-induced ion drift can have a significant effect on the magnetic field. In principle, the initial magnetic field may be enhanced (or weakened). In this connection it is of interest to consider the following possible mechanism for the generation of magnetic field in the atmospheres of magnetic stars.

It is obvious that the velocity of ion drift in the stellar atmosphere induced by stellar radiation possesses only the radial component u_k in the absence of magnetic field. If the stellar dipole magnetic field is taken into account, the ions drifting in the atmosphere are subjected to an averaged Lorenz force, which is directed perpendicular to the radial velocity u_k and the magnetic field B and produces circular electric currents around the star. These currents can either enhance the initial stellar magnetic field, or weaken it, depending on the direction of the radial velocity of the ion drift (outside or inside the star).

The lowering of ion drift velocity with increasing magnetic field (for $\omega_c \gtrsim \nu_\nu$) exerts a stabilising effect on the field build-up by this mechanism and is responsible (if this mechanism of field generation is the principal one) for the field saturation at a level defined, by the order of magnitude, by the condition $\omega_c/\nu_n \sim 1$. This gives $B \sim Mc\nu_n/e$, and therefore for the ion mass $M \sim 40$ amu and a transport ion collision frequency $\nu_n \sim 10^7$ s^{-1} typical of the atmospheres of magnetic stars, we obtain the estimate

$B \approx 4 \times 10^4$ G. This corresponds to the strongest magnetic fields detected in the atmospheres of magnetic stars [14–16].

Note that the atmospheres of magnetic stars are stable (not subject to convection) [15, 16]. The atmospheric stability is a necessary condition both for the mechanism of chemical element separation and for the mechanism of magnetic field generation due to light-induced ion drift.

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