favorable because of the quantization of the electron motion in the transverse direction. If $\hbar^2/m^*d^2 > V_0 - \Delta$, then all the quasiparticles produced under the influence of the electrostatic image forces populate only the lower level of the transverse motion. We then return to the already discussed case of a two-dimensional semiconductor.

The realignment of the ground state of a semiconductor under the influence of electrostatic image forces can take place also for extrinsic semiconductors (with impurity concentration $N \ll 1/d^3$ and with a ground-state radius $\rho_0 \ll d$). The electrostatic-image forces decrease the impurity ionization energy and accordingly increase the radius of the impurity state[5]. If the impurity-band broadening due to the overlap of the impurity orbitals were to become comparable with the decreased ionization energy, a semiconductor - metal transition would take place.

We note that in the case of extrinsic semiconductors with a large width of the forbidden band (and hence with small ε) and with a donor or acceptor depth $\Delta \sim 10^{-2}$ eV, the electrostatic image forces become appreciable even for layers several hundred Angstrom thick. This circumstance would facilitate the observation of the realignment in question.

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ACCELERATION OF ATOMS BY A STATIONARY FIELD

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This paper deals with a special resonant-field configuration in which the atom is acted upon by a spatially constant field. The excited atom can be accelerated in this case to an energy of 1 keV. This effect can be used for spatial separation of excited and non-excited atoms in an atomic beam.

In a strong resonant field, the atom is acted upon by an appreciable force, on the order of $10^2 - 10^3$ eV/cm. This estimate follows from the general formula for the force F acting on the dipole moment of the atom p in an inhomogeneous field E

$$F = p\nabla\hat{E}^* + \text{c.c.} \tag{1}$$

If we put $p \sim 1$ Debye, $k \sim 10^{-5}$ cm^{-1} for the wave number, and $E \sim 3 \times 10^5 - 3 \times 10^6$ V/cm. Formula (1) has been written out in the resonance approximation, and the frequencies of the field and of the dipole moment are reckoned from the frequency of the working transition.

In a quasistationary field, the induced dipole moment follows the field adiabatically, and the force (1) takes the form

$$F = \nabla \langle \hat{H} \rangle \tag{2}$$

[5] Similar effects can probably take place also for adsorbed atoms.

where $\hat{H} = \hat{d}\hat{E}^{*} + \text{h.c.}$ is the Hamiltonian of the electron interaction with the resonant field, and $\langle d \rangle = p$.

Under ordinary conditions this force does not come into play, since it oscillates over the wavelength. If, however, the field is taken in the form of a uniformly-accelerated sinusoid, then the captured atoms will have an acceleration proportional to the rate of change $\dot{\omega}$ of the field frequency [1]. At sufficiently large $\dot{\omega}$, the acceleration effect is such as if the force F were to be constant in space. Thus, for example, a hydrogen atom can be accelerated to 1 keV within 10^{-7} sec.

The force (1) acts only on excited resonant atoms. It can therefore be used for spatial separation of excited atoms from atoms in the ground state, just as in an ammonia maser. One can thus obtain inverted population with respect to the ground state and generation in the violet or ultraviolet region of the spectrum.

To effect such a separation, it suffices to impart to the excited atom in the beam a transverse momentum of the order of the thermal momentum. From this point of view it would be useful to study other acceleration mechanisms not connected with pulsed acceleration of the captured atoms.

Ashkin [2] proposed to use the pressure of light for resonant acceleration of atoms. In this case, the atom acquires during the lifetime the momentum of the resonant photon. The ponderomotive force (1) exceeds the light-pressure force by a factor $dE\tau/\hbar$, where τ is the lifetime of the atom. This parameter can be very large in a strong field.

It is of interest to consider the possibility of accelerating the atoms with a stationary field. It is clear that this is possible only if the force is not in the form [2], i.e., the adiabatic approximation is violated. We consider below a simple example of such a field configuration, in which the effective force has a definite sign over a distance much larger than the wavelength, and we estimate the possible acceleration effect.

In a field of the form

$$E = E_0(x) + E_1(x, y)e^{i\omega_1 t}; \quad E_1 \ll E_0 \qquad (3)$$
$$E_0(x) = E_0 \cos kx$$

the work is performed mainly by the strong resonant field

$$F = 2p_0 \nabla E_0(x) \qquad (4)$$

where p_0 is the value of the dipole moment at the initial instant of time. Using

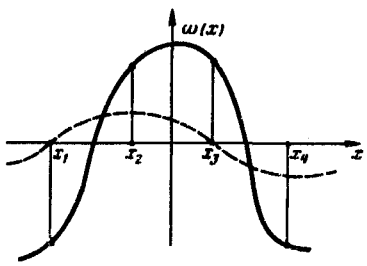

Fig. 1

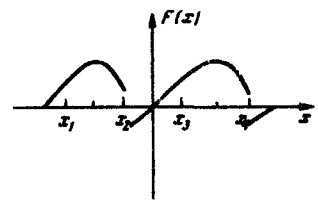

Fig. 2

the analogy with the magnetic moment, it can be said that p_0 is a conserved component directed along the effective magnetic field, the role of which is played by the quantity $E_0(x)$. The damping p_0 due to the spontaneous emission can be neglected if the damping time is sufficiently long, on the order of 10^{-7} sec.

The atomic-precession frequency $\omega(x) = 2dE_0(x)$, where d is the matrix element of the dipole moment of the transition, varies with the coordinate and its value is of the order of $10^{12} - 10^{13}$ Hz. The Doppler frequency is $kv_0 \sim 10^9$ Hz (v_0 is the initial velocity of the atom).

The time required to negotiate a distance equal to half the wavelength, over which the force (4) has a constant sign, is thus larger by 3 - 4 orders of magnitude than the precession period. It is therefore possible to reverse the sign of a conserved component of the dipole moment p_0 approximately at the place where the sign of $\nabla E_0(x)$ is reversed. This can be realized with the aid of a weak resonant field $E_1(x, y)$, which we choose in the form

$$E_1(x, y) = E_1 e^{i\Delta k y} \cos(kx + \pi/4) \tag{5}$$

$$\omega_1 = \omega(x_2), \quad kx_2 = -\pi/4 \tag{6}$$

The resonance condition can be readily understood with the aid of Fig. 1, in which the solid line shows the local precession frequency $\omega(x)$ and the dashed line shows $|E_1(x, y)|$.

At the point $x = x_2$ the atom goes through the resonance (6). If the condition

$$dE_1 \gg \sqrt{\hbar k v dE_0} \tag{7}$$

is satisfied at the same time, then we deal with a slow passage through resonance. Then p_0 reverses sign on passing through the point x_2 regardless of the phase of the field and of the oscillator [2 - 4].

At $x = x_3$, the precession frequency $\omega(x_3)$ and the signal frequency are also equal, but the amplitude E_1 vanishes. We then deal with rapid passage through resonance and p_0 remains unchanged. The next reversal of the sign of p_0 occurs only at $x = x_4$.

The force is shown schematically in Fig. 2. At the points x_2 and x_4 the force experiences a discontinuity, but actually the width of the transition region is of the order of $2\pi E_1/kE_0$. Thus, after traversing a wavelength the atom energy changes by an amount

$$U = 8p_0 E_0 \cos \pi/4. \tag{8}$$

We now estimate the number N of the traversed wavelengths, at which the assumptions made above are violated. The limitation on N is caused by the presence of small changes of p_0, which depend on the phase of the field and of the atom, on passing through resonance. In the case of rapid passage through resonance these corrections are of the order of E_1/E_0. Taking into account the random character of these corrections, we can estimate the number of steps in which p_0 changes by an amount on the order of unity at $N = (E_0/E_1)^2$.

In the case of slow passage through resonance, the corrections that depend on the random phases are of the order of $\exp(-\sqrt{v_{max}/v})$, where $v_{max} = dE_1^2/kE_0$. In view of the strong dependence on the velocity, we can assume approximately that the random corrections to p_0 become significant at $v \approx v_{max}$. On the other hand, in the case of appreciable acceleration we have $v = \sqrt{2NU/M}$, where M is

the mass of the atom. This yields the optimal value of the signal

$$E_1 \approx E_0\left[\frac{(\hbar k)^2}{2MdE_0}\right]^{1/16} \tag{9}$$

at which the atom traverses the maximum number of wavelengths as it accelerates

$$N_{max} \approx \left[\frac{2MdE_0}{(\hbar k)^2}\right]^{1/3}. \tag{10}$$

We see thus that the acceleration is more effective for heavy particles with the aid of a long-wave field. The final energy NU of the atom does not depend on the initial velocity. It is convenient, however, to express (10) also in the form

$$N_{max} \approx \left(\frac{dE_0}{\hbar k v_0}\right)^{2/3} \eta^{-1/3}, \tag{11}$$

where $\eta = 2U/Mv_0^2$.

In a strong field $U \sim 10^{-2}$ eV, so that $\eta^{1/3} \approx 1$ at energies corresponding to room temperatures. $dE_0/\hbar k v_0 \sim 10^3$, so that $N_{max} = 2U/Mv_0^2$.

In place of a field E_1 that changes the dipole phase by π, we can use a series of π pulses tuned to the maximum precession frequency $\omega(x)$. In this case those atoms which pass through an antinode of the strong field also change the phase of the dipole by π in the presence of a short-duration π-pulse, and their energy changes by dE after traversing one wavelength. The acceleration has then a stochastic character. Even if we assume that the atom can change its energy by dE only once, this may be sufficient to permit the excited atom to leave the atomic beam.

Finally, when the atom crosses a narrow light beam whose intensity is time-modulated, its energy can also change by an amount on the order of dE. To this end it suffices to have the period of the modulation of the order of the time of flight.

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