

In the case of low injection levels at the same temperatures (30°), the relative role of the radiation connected with electron capture by Zn^- increases, and the intensity of the emission produced when the electrons go over into the valence band decreases (curve III). Thus, apparently no phononless transitions of electrons to Zn^0 take place.

The activation energy determined from the position of the emission line for Zn^0 turns out to be 0.03 eV, in good agreement with the data obtained by other methods [6]. It is obvious that the absence of a barrier in the capture of an electron by Zn^0 eliminates the difference between the thermal and optical ionization energies.

It is interesting to note that optical transitions of electrons to different charge states of Zn in Ge appear in sequence, depending on the temperature interval, in investigations of one and the same sample. At low temperatures, when the zinc atoms are in the $Zn^{=}$ charge state, we obviously have a consecutive capture of two holes, so that the zinc atom goes over into the Zn^0 state ($Zn^{=} + p + p \rightarrow Zn^0$). This is followed by radiative capture of an electron by the Zn^0 ion. At higher temperatures, owing to the thermal ejection of holes from the 0.03 eV level, transitions predominantly to the 0.09 eV level appear.

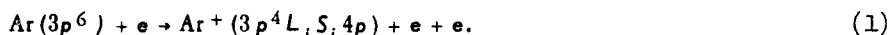
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- [1] Ya. E. Pokrovskii and K. I. Svistunova, FTT 6, 19 (1964), Soviet Phys. Solid State 6, 13 (1964).
- [2] C. Benoit a la Gnillaune and J. Cernogora, Symp. on Radiative Recombination in Semiconductors, Paris, 1964, p. 121.
- [3] L. V. Keldysh, JETP 45, 364 (1963), Soviet Phys. JETP 18, 253 (1964).
- [4] V. L. Bonch-Bruevich, FTT 4, 298 (1962) and 5, 2890 (1963), Soviet Phys. Solid State 4, 215 (1962) and 5, 2117 (1964).
- [5] M. Zax and J. Hopfield, Phys. Rev. 124, 115 (1961).
- [6] V. I. Sidorov, FTT 5, 3006 (1963), Soviet Phys. Solid State 5, 2199 (1964).

IONIZATION OF Ar ATOMS IN THE Ar^+ EXCITED STATE

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Recent publications point to an important role played by electron-atom collision processes in which the states of two atomic electrons are changed. An example of such a double process is the ionization of an Ar atom with simultaneous excitation of Ar^+ [1,2]:



When the atom is described by orthogonal single-electron wave functions, the cross section of this process in the first Born approximation is equal to zero. The sudden-perturbation ap-

proximation is used in [1,3] to estimate the cross section. It is assumed that the knocked-out electron leaves the atom so rapidly that the state of the remaining atom has not time to change. The ratio of the cross sections of the double and single processes at high energies is equal to Δ^2 , where Δ is the non-orthogonality integral of the wave functions of the atom and ion. Qualitatively similar results are obtained in the first Born approximation if the non-orthogonality of the atomic functions in the initial and final states is taken into account [4]. Here, however, the change of the wave function of the knocked-out electron upon excitation of the ion, and the difference between the thresholds, are taken into account in a more natural manner.

Since Δ is small, the first-order cross section σ^I is also small, and an important role may be assumed by the second order of perturbation theory, in which (1) is regarded as a transition through the intermediate levels $3p^5 4p$ and $3p^5 EI$ (E and I are the energy and angular momentum of the knocked-out electron). σ^{II} does not depend on Δ .

The total cross section of the process (1) is equal to

$$\sigma = \frac{8}{k_0^2} \sum_{l_0}^{\epsilon} \int_{k_0-k_1}^{\epsilon} \frac{k_0+k_1 dq}{q^3} \sum_{L,S,M_L,M_S} |F(q, k_1)|^2, \quad q = k_1 - k_0, \quad (2)$$

where ϵ is the threshold energy, and $L, S, M_L,$ and M_S are the total angular momenta of the system $Ar^+ + e$ (the subscripts 0 and 1 refer to the initial and final states).*

The scattering amplitude is $F = F^I + \sum_a F_a^{II}$ (summation over the intermediate states). Estimates have shown that the main contribution to $\sum_a F_a^{II}$ is made by $a = 3p^5 EI$. If we neglect furthermore the interference between F^I and F^{II} , then $\sigma = \sigma^I + \sigma^{II}$. We calculated σ^I in [4]. F_a^{II} can be written in the form

$$F_a^{II}(q, k_1) = \int dR e^{-iqR} U_{00}(R) \int dx U_{01}(R-x) \frac{e^{ik_0 x + ik_1 x}}{2\pi x}. \quad (3)$$

$$U = \sum_j |\mathbf{R} - \mathbf{r}_j|^{-1}, \quad k_0^2 = k_1^2 + 2(E_0 - E_\sigma) = k_1^2 + 2(E_1 - E_\sigma).$$

When $k_1 = 0$ formula (3) simplifies greatly [5]:

$$F_a^{II}(q) = \frac{2}{K^2} \int dR e^{-iqR} U_{00}(R) Q_{01}(R) dR \quad (4)$$

$$Q = \sum_j [\exp\{i\sqrt{2(E_1 - E_\sigma)}|\mathbf{r}_j - \mathbf{R}|\} - 1] |\mathbf{R} - \mathbf{r}_j|^{-1} \quad (5)$$

$$K^2 = 2(E_1 - E_\sigma) \equiv k_{01}^2. \quad (6)$$

On the other hand, using the quasiclassical approximation, we can show that at high energies $F^{II} \sim k_a^{-1}$. Since the practical utilization of (3) entails great difficulty, we shall use an approximate interpolation formula. Namely, we set F_a^{II} equal to (4), but in lieu of (6) we assume that

$$K^2 = k_\sigma k_{\sigma 1}. \quad (7)$$

This approximation was verified in the case of double excitation of Ca, where an exact calculation of F^{II} is possible [5].

For large $(E_1 - E_a)$, we can neglect the rapidly oscillating quantity in (6), and then (6) coincides (accurate to the definition of the factor K) with the known adiabatic approximation, in which the Green's function $[\exp(ik_a x)]/2\pi x$ is replaced by $-2\delta(x)/k_a^2$. Thus, the condition for the applicability of the adiabatic approximation, besides smallness of k_1 , is that the distance from the final to the intermediate level be large.

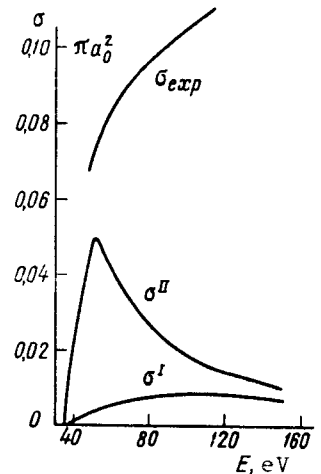
Summing over the momenta (in the LS-coupling scheme), integrating over the angles, and confining ourselves for simplicity in the multipole expansion of Q_{a1} , we get

$$\sum_{LSM_L M_S} |F_{II}(q)|^2 = 30 \left\{ \frac{q^2}{K^2} G_{L_i S_i}^{3P} \right\}^2 \sum_{\lambda=l\pm 1} \frac{2(l+\lambda+1)}{2\lambda+1} \left| \int_0^\infty y_{00}^\lambda Z_{01i\lambda}(qr) r^2 dr \right|^2 \quad (8)$$

$$y_{00}^\lambda(r) = \int \frac{r^\lambda}{r^{\lambda+1}} R_{3p}(r') R_{EI}(r') r'^2 dr'$$

$$Z_{01}(r) = \int [i k_{01} i_0(k_{01} r_<) h_0^{(1)}(k_{01} r_>) - \frac{1}{r_>}] R_{3p}(r') R_{4p}(r') r'^2 dr', \quad (9)$$

where G is the Racah fractional-parentage coefficient, R_{3p} and R_{4p} are the semi-empirical radial functions of $Ar(3p^6)$ and $Ar^+(3p^4 4p)$, R_{EI} is the radial function of the knocked-out electron in the field of $Ar^+(3p^4 4p)$, and j_λ and $h_0^{(1)}$ are spherical Bessel and Hankel functions (the latter is complex, but this affects the calculations very little). The sum over I in (2) was taken up to I = 6. The ratios of the effective excitation cross sections of the terms $L_i S_i$ of the initial ion Ar^{++} are proportional to G^2 for σ^{II} and to $G^2 \Delta^2$ for σ^I . The values of Δ^2 calculated with the aid of semi-empirical wave functions [6] are: $\Delta^2(3p) = 1.4 \times 10^{-3}$, $\Delta^2(1s) = 3.5 \times 10^{-3}$, and $\Delta^2(1D) = 4.5 \times 10^{-3}$. The figure shows plots of σ^I and σ^{II} for the $3p$ term of Ar^{++} . It must be noted, however, that Δ , and consequently also σ^I , is quite sensitive to the level energy. In particular, violation of the LS-coupling scheme, which is significant in the case of Ar, leads to a mixing of terms with different energies, i.e., to an appreciable change in the "effective value" of Δ . Estimates show that apparently in this case the cross section σ^I summed over LS increases considerably, whereas σ^{II} remains small. **



The same figure shows an experimental curve from a recent paper [7] (the cross section summed over seven levels). The experimental cross sections exceed the theoretical ones, probably as a result of violation of the LS coupling and owing to the role of the cascades (the latter is noted by the authors of [7]).

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[1] W. R. Bennet, Appl. Opt. Suppl. 2, 3 (1965).

- [2] S. Koozekanani, Quantum Electronics QE-2, 87, IB⁴ (1966).
 [3] M. H. Mittleman, Phys. Rev. Lett. 16, 498 (1966).
 [4] L. A. Vainshtein and A. V. Vinogradov, Optika i spektroskopiya 23 (1967), in press.
 [5] I. L. Beigman, L. A. Vainshtein, and A. V. Vinogradov, FIAN Preprint No. 5, 1967.
 [6] L. A. Vainshtein, Optika i spektroskopiya 3, 313 (1957).
 [7] W. R. Bennet, Jr., G. N. Mercer, P. I. Kindlmann, B. Wexler, and H. Hyman, Phys. Rev. Lett. 17, 987 (1966).

* We use atomic units with πa_0^2 equal to unity for σ .

** According to (8), only the ²P term is excited in the LS coupling scheme.

SELF-ACTION OF LIGHT IN CRYSTALS

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1. The possibility of observing self-focusing [1-5] and other "self-action" effects of light [6] is presently connected almost exclusively with the cubic nonlinearity of the dependence of the polarization on the field in liquids, since the latter are characterized by large values of the increment to the nonlinear polarizability [7,8]. For crystals, the corresponding values are usually much lower. In crystals without inversion center, however, strong self-action is possible as a result of the second-harmonic reaction. In this case the effects due to anisotropy are distinctive and of interest. We consider below the peculiarities of such a mechanism and present very simple estimates.

The interaction of waves having frequencies ω and 2ω in a quadratic medium leads to the appearance, at the corresponding frequencies, of the nonlinear-polarization vectors

$$P_i^{(2\omega)} = \chi_{ikl}^{(2\omega)} E_k^{(\omega)} E_l^{(2\omega)}(a); P_i^{(\omega)} = \chi_{ikl}^{(2\omega-\omega)} E_k^*(\omega) E_l^{(2\omega)}(b). \quad (1)$$

By determining $\vec{E}^{(2\omega)}$ from (1a) in the given field approximation (which is valid everywhere except in the synchronism directions), and substituting the result in (1b), we obtain for $P_i^{(\omega)}$ in a quasi-plane wave an expression of the form

$$P_i^{(\omega)} = \theta_{iklm} A_k^* A_l A_m e^{-ik \cdot r} + P_i' e^{-i(k_1 - k_2) \cdot r}. \quad (2)$$

Here \vec{A} is the amplitude of the fundamental wave, and $\vec{k}_{1,2}$ are the values of the wave vector at the frequencies ω and 2ω respectively; $\hat{\theta}$ is a tensor proportional to the product of the tensors $\hat{\chi}^{(2\omega)}$ and $\hat{\chi}^{(2\omega-\omega)}$; P_i' is the polarization connected with the proper waves at the frequency 2ω and determined from the boundary conditions. Inasmuch as the crystal dimensions are usually much larger than the coherence length $l_{\text{coh}} \sim \pi/|2\vec{k}_1 - \vec{k}_2|$, the term P_i' can be