Calculation of transition probabilities and lifetime for He atom Stark states in the alternating electric field

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Received November 18, 2005

The dependence of transition probabilities on the frequency and the strength of a circularly polarized electric field is studied theoretically for the He atom. Regularities in the behavior of the probabilities were revealed and investigated. It was shown that the transition probabilities and lifetimes of the Stark energy levels have the polynomial dependence on the electric field strength. It was found that an increase in the frequency of the electric field leads to a decrease in the sensitivity of the transition probabilities to the electric field strength. An anisotropy of the probabilities for transitions between Stark levels was studied.

Introduction

The gas discharge electric field and transition probabilities are among the most important discharge characteristics. In plasma physics, the Stark effect is used widely in plasma diagnostics, particularly, to determine the energy distribution function, electron temperature, electron density, etc. Transition probabilities are necessary to calculate the lifetimes of energy levels and intensities of spectral lines of atoms in the external electric field. Calculation of the Stark effect and the probabilities of transitions between Stark energy levels is an especially urgent problem for noble gases, because they are widely used in plasma physics. For reliable interpretation of experimental results, they, obviously, should be compared with theoretical calculations of the corresponding characteristics. Consequently, a reliable theoretical method, determining wave functions for an atom in the external electric field, is needed. The wave functions are used to calculate the Stark effect and any characteristics of the atom behavior in the electric field.

Spectra of atoms subject to the alternating electric field action can be determined from the nonstationary Schrödinger equation. The methods of this equation solution depend on the field polarization (linear, circular, or elliptical). This paper considers the dynamic Stark effect in the circularly polarized electric field. The electric fields of this polarization arise in the high-frequency discharge[1] and upon the laser excitation.[2] At the field circular polarization, due to separation of spatial and temporal variables, the Schrödinger equation solution becomes much simpler, namely, the nonstationary Schrödinger equation reduces to the stationary one in the rotating wave approximation.[3] In spite of the simplification, the stationary Schrödinger equation solution is still a complicated problem, since the perturbation theory can be applied only under the following restrictions. First, the electric field strength should be relatively low and the perturbation, induced by the external electric field, should be smaller than the separation between neighboring energy levels. Second, the resonance and nonresonance perturbations should be calculated by different methods. Finally, to calculate the excitation of an atom by the low-frequency or high-frequency fields, it is also necessary to apply different methods.[4–6]

Attempts to derive equations for calculating shifts and splitting of atomic energy levels in the field of an arbitrary frequency and strength were repeatedly undertaken earlier. However, these attempts turned out successful only in some particular cases. In Ref. 7, the equations are derived for calculation of energy level shifts for different model systems; the calculations for systems in single- and two-level approximations have been carried out in Refs. 6 and 8. In Refs. 9 and 10, such equations were obtained for a particle in a short-range potential and for a negative ion. General equations for calculation of the Stark effect in terms of the perturbation theory were derived in Ref. 11. However, they can be applied only to calculations of an isolated atomic level in the absence of resonance with the field. In addition, if atoms with shells including equivalent electrons are considered, then the equations presented in Ref. 11 can be applied only to calculation of Rydberg states of these atoms.

In this paper, to obtain the wave functions for atoms in the circularly polarized electric field, the theoretical method proposed and developed in Refs. 12 and 13 is used. This method, free of restrictions inherent in the perturbation theory, is used to calculate the dynamic Stark effect and the probabilities of spontaneous transitions between the Stark levels for the He atom, as well as to study the dependence of the Stark effect and the transition probabilities on the frequency and strength of the electric field. This problem is of urgency, because both the Stark effect and transition probabilities are widely used for explanation of the processes proceeding in plasma and for the plasma diagnostics.

0235-6880/06/07 519-07 $02.00

Method of calculation

In the circularly polarized electric field, the nonstationary Schrödinger equation is written in the form

$$i\frac{\partial\psi_n(\mathbf{r},t)}{\partial t}=\left[\hat{H}_0(\mathbf{r})-eF(x\cos\omega t\pm y\sin\omega t)\right]\psi_n(\mathbf{r},t), \quad (1)$$

where ψ_n is the wave function of the system nth state; $\hat{H}_0(\mathbf{r})$ is the unperturbed Hamiltonian, the operator $-eF(x\cos\omega t\pm y\sin\omega t)$ describes the perturbation caused by the atom interaction with the circularly polarized field of the frequency ω and the strength F. The signs + and − correspond to the right and left polarizations of the field, respectively. To pass on to the stationary Schrödinger equation, we use the rotating wave approximation.[3]

To pass on to the coordinate system rotating around the axis Z with the frequency ω, we introduce the wave function in this coordinate system

$$\varphi(\mathbf{r},t)=\exp(i\omega t\hat{J}_z)\psi(\mathbf{r},t), \quad (2)$$

where $\hat{J}_z$ is the z-component of the total angular momentum operator. After substitution of the wave function (2) into Eq. (1), we have

$$i\frac{\partial\varphi(\mathbf{r},t)}{\partial t}=\hat{Q}\varphi(\mathbf{r},t),\quad \hat{Q}=(\hat{H}_0-\omega\hat{J}_z\pm F\hat{D}_x). \quad (3)$$

As can be seen from Eq. (3), the operator $\hat{Q}$ is independent of time. Consequently, in the rotating wave approximation, it is possible to pass on from the nonstationary Schrödinger equation (1) to the stationary one, and we obtain

$$\hat{Q}\varphi(\mathbf{r})=\varepsilon\varphi(\mathbf{r}), \quad (4)$$

where

$$\varphi(\mathbf{r},t)=\exp(-i\varepsilon t)\varphi(\mathbf{r}); \quad (5)$$

$\hat{Q}$ is the energy operator of the atom in the electric field; ε and $\varphi(\mathbf{r},t)$ are the energy and the wave function of the atom in the electric field in the rotating coordinate system. Obviously, ε and $\varphi(\mathbf{r},t)$ can be found through the use of the stationary perturbation theory. Instead of the Schrödinger equation (4) solution within the framework of the perturbation theory, we use the approach proposed in Refs. 12 and 13.

It was shown in Ref. 12 that the wave functions and energies of an atom, being the solutions of the Schrödinger equation (4), are determined by diagonalization of $\hat{Q}$. This matrix can be obtained in representation of the unperturbed wave functions $\varphi_n^{(0)}$, calculated in the absence of the external electric field. In this representation, the matrix elements of $\hat{Q}$ are written as

$$Q_{mn}=E_n^{(0)}\delta_{mn}-\omega<\varphi_m^{(0)}(\mathbf{r})\left|\hat{J}_z\right|\varphi_n^{(0)}(\mathbf{r})>\pm F<\varphi_m^{(0)}\left|\hat{D}_x\right|\varphi_n^{(0)}>, \quad (6)$$

where $E_n^{(0)}$ is the energy of the nth state of the atom in the absence of the external electric field; D_x is the x-component of the dipole transition operator.

The diagonalization of the energy matrix $\hat{Q}$ with the elements (6) yields a set of wave functions and the energy spectrum for n states of the atom in the electric field. Diagonalization of $\hat{Q}$ allows us to obtain the energies ε_n and the wave functions in the form

$$\varphi_n(\mathbf{r},t)=\mathrm{e}^{-i\varepsilon_n t}\sum_k C_{nk}\varphi_k^{(0)}(r) \quad (7)$$

for n states of the atom in the external electric field in the rotating coordinate system. The coefficients C_{nk} in the wave function (7) depend on the frequency and strength of the external electric field. To find the atom averaged energies in the initial coordinate system, it is necessary to carry out the averaging over the period of oscillations. After the averaging, the energy of the system in the electric field in the initial coordinate system is written in the following form:

$$\bar{E}_n=<\psi_n(\mathbf{r},t)\left|H(\mathbf{r},t)\right|\psi_n(\mathbf{r},t)>=$$
$$=\varepsilon_n+\omega<\varphi_n(\mathbf{r})\left|\hat{J}_z\right|\varphi_n(\mathbf{r})>. \quad (8)$$

It follows from Eq. (8) that $\bar{E}_n$ is independent of time.

The matrix elements of the operator D_x are determined as

$$<\varphi_m^{(0)}\left|\hat{D}_x\right|\varphi_n^{(0)}>=<\gamma JM|D_x|\gamma'J'M'>=\frac{(-1)^{J-M}}{\sqrt{2}}\times$$
$$\times\left[\begin{pmatrix}J & 1 & J'\\ -M & -1 & M'\end{pmatrix}-\begin{pmatrix}J & 1 & J'\\ -M & 1 & M'\end{pmatrix}\right]<\gamma J\|D\|\gamma'J'>, \quad (9)$$

where the reduced matrix elements $<\gamma J\|D\|\gamma'J'>$ are calculated depending on the bond type. For the He atom, the LS bond takes place, and the matrix elements $<\gamma J\|D\|\gamma'J'>$ for the transitions $l_1^{N_1}l_2-l_1^{N_1}l_3$ are calculated by the equations[14]:

$$<\gamma J\|D\|\gamma'J'>=$$
$$=<l_1^{N_1}l_2\alpha_1L_1S_1l_2S_2LSJ\|D\|l_1^{N_1}l_3\alpha_1'L_1'S_1'l_3S_3L'S'J'>=$$
$$=(-1)^{\varphi}\delta(\alpha_1L_1S_1,\alpha_1'L_1'S_1')Q(\alpha,LS,L'S',JJ')\times$$
$$\times<n_2l_2\|r\|n_3l_3>, \quad (10)$$

where

$$\varphi=L_{\mathrm{core}}+S+l_3+J';$$
$$Q(\alpha,LS,L'S',JJ')=$$
$$=\delta(S,S')\sqrt{(2L+1)(2L'+1)(2J+1)(2J'+1)}\times$$
$$\times\begin{Bmatrix}L & J & S\\ J' & L' & 1\end{Bmatrix}\begin{Bmatrix}l_2 & L & L_{\mathrm{core}}\\ L' & l_3 & 1\end{Bmatrix}; \quad (11)$$

$$< n_2 l_2 \| r \| n_3 l_3 >=$$

$$= (-1)^{\frac{l_2+l_3+1}{2}+l_2} \sqrt{\max(l_2, l_3)} < R_{n_2 l_2} | r | R_{n_3 l_3} >,$$

$$< R_{nl} | r | R_{n'l'} >= \int_0^\infty R_{nl}(r) r R_{n'l'} r^2 \mathrm{d}r; \qquad (12)$$

L_{core} is the orbital quantum number of the atomic core without the outer l_2 electron. The radial integral $< R_{nl} | r | R_{n'l'} >$ from Eq. (12) is calculated by the semiempirical equation, being the improved modification of the Bates–Damgaard formula.[15] The particular form of the semiempirical equation and the details of its derivation can be found in Ref. 16.

The wave functions and energies, determined from the diagonalization of the $\hat{Q}$ matrix, are used to calculate the probabilities of spontaneous transitions in the electric field. The probability of spontaneous emission of a photon into an element of the solid angle dΩ upon the transition from the state $|n>$ to the state $|m>$ with the polarization $\mathbf{e}_q$ is determined by the equation[14]:

$$A_q = \frac{e^2 \omega}{hc^3 m^2} \left| \mathbf{e}_q < \Psi_n | \mathbf{p} \mathrm{e}^{i\mathbf{kr}} | \Psi_m > \right|^2 \mathrm{d}\Omega. \qquad (13)$$

Here ω is the transition frequency; $\mathbf{p}$ is the electron momentum; $\mathbf{e}_q$ is the vector of polarization; $\mathbf{k}$ is the wave vector. In the dipole approximation, Eq. (13) can be rewritten as

$$A_q = \frac{\omega^3}{hc^3} \left| \mathbf{e}_q < \Psi_n | \mathbf{D} | \Psi_m > \right|^2 \mathrm{d}\Omega, \qquad (14)$$

where $\mathbf{D} = -e \sum_i \mathbf{r}_i$ is the dipole moment of the atom; Ψ_n and Ψ_m are the wave functions of the nth and mth states of the atom in the external electric field. On the basis of Eq. (14), the total transition probability for radiation, polarized in the direction $\mathbf{e}_q$ and averaged over all possible spatial orientations in $\mathbf{D}$, is calculated by the equation

$$A_{nm} = \frac{4\omega^3}{3\hbar c^3} \sum_q \left| < \Psi_n | D_q | \Psi_m > \right|^2, \qquad (15)$$

where D_q are cyclic components of $\mathbf{D}$. The wave functions Ψ_n and Ψ_m are determined from the diagonalization of the energy matrix $\hat{Q}$ with the matrix elements (6). After substitution of Ψ_n and Ψ_m into Eq. (15) and application of the Wigner–Eckart theorem, the equation for the probability of the transition $JM \to J'M'$ between the magnetic sublevels takes the form

$$A(JM \to J'M') = \frac{4\omega^3_{JM,J'M'}}{3\hbar c^3} | D_{JM,J'M'} |^2;$$

$$D_{JM,J'M'} =$$

$$= \sum_q \sum_{ij} C_i^{(JM)*} C_j^{(J'M')} (-1)^{J_i - M_i} \begin{pmatrix} J_i & 1 & J_j \\ -M_i & q & M_j \end{pmatrix} \times$$

$$\times < \gamma_i J_i \| D \| \gamma_j J_j >, \qquad (16)$$

where $C_i^{(JM)}$ and $C_j^{(J'M')}$ are the coefficients of expansion of the wave functions for the atom in the filed in terms of the unperturbed wave functions $\varphi_i^{(0)}(\gamma_i J_i M_i)$; $\omega_{JM,J'M'}$ is the frequency of transition $JM \to J'M'$. The reduced matrix elements $< \gamma_i J_i \| D \| \gamma_j J_j >$ are calculated by Eqs. (10)–(12).

The existing experimental techniques usually do not permit the measurement of the split between magnetic sublevels and the determination of the corresponding transition probabilities. That is why we calculate the probabilities of transitions $J \to J'$ between the Stark energy levels, using the equation

$$A(J \to J') = \frac{1}{2J+1} \sum_{MM'} A(JM \to J'M'). \qquad (17)$$

The lifetime of the Jth energy level in the electric field is determined as

$$\tau_J = \frac{1}{A_J} = \frac{1}{\sum_{J'} A(J \to J')}. \qquad (18)$$

Consider the results obtained within the framework of the above approach.

Results and discussion

The theoretical method proposed was applied to calculation of the dynamic Stark effect and the transition probabilities for the He atom in the circularly polarized electric field. All calculations were performed in the LS bond. When calculating the $\hat{Q}$ matrix, the ns, np, nd, and nf states with $n \le 10$ were taken into account. Thus, the calculation of the helium atom energy matrix in the electric field took into account 115 energy levels (501 magnetic sublevels). The electric field strength was considered in a range up to 1 kV/cm with the following frequencies ω: 100, $241.813 \cdot 10^5$, and $283.005 \cdot 10^6$ MHz. Such electric fields are produced by real excitation sources: ω = 100 MHz is the electric field frequency in a high-frequency discharge, ω = $241.813 \cdot 10^5$ MHz is the frequency of the CO_2 laser, and ω = $283.005 \cdot 10^6$ MHz is the frequency of the Nd laser. In this work, spectral lines of the visible region were considered.

First of all, the calculations have shown that the increase in the frequency of the electric field results in the decrease of the shift and split of the atomic energy levels in the field. As an illustration, Fig. 1*a* depicts the spectral line shift and split for the He atom in the electric field of different frequencies. The

shift is determined as $\Delta E = E - E_0$, where E is the position of the line under consideration in the electric field, E_0 is the position of this line in the field absence.

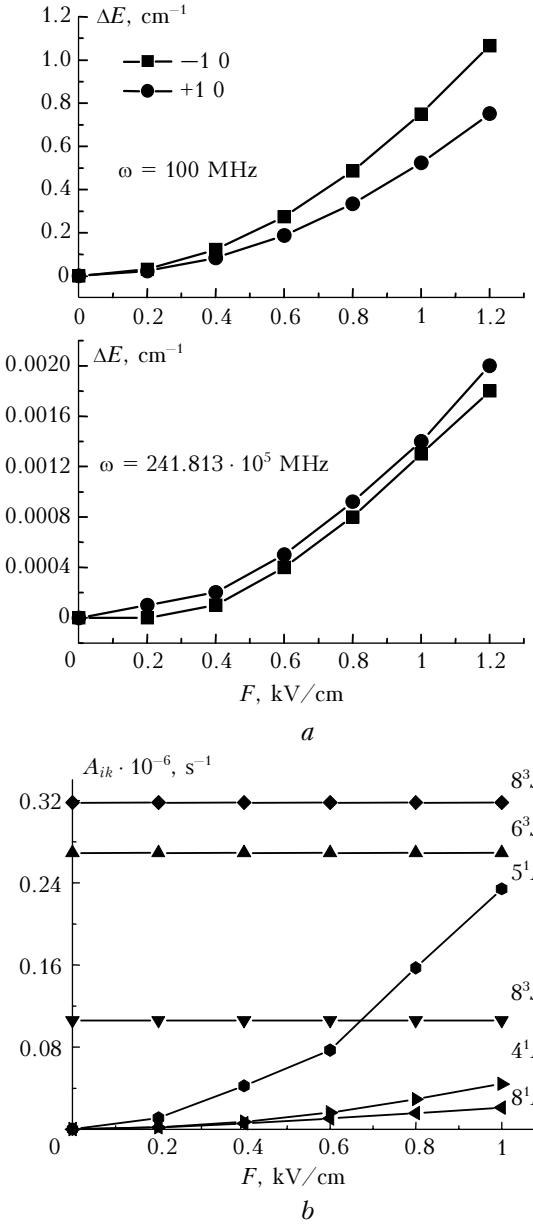

Fig. 1. Stark effect and transition probabilities as functions of the electric field strength for the He atom: 8^1P_1–2^1S_0 spectral line, λ = 329.772 nm (*a*); probabilities of allowed and forbidden transitions (*b*).

The analysis of the wave functions has shown that the increase in the electric field frequency leads to the decrease in interaction between the atom energy levels in the electric field, which is one of the possible causes for the decrease in the spectral line shift and split with the increasing electric field frequency. In the high-frequency discharge, the interaction between energy levels is very strong, which leads to appearance of forbidden lines, and probabilities of forbidden transitions become comparable with those of allowed transitions, as the electric field strength increases (Fig. 1*b*). With the increase in the electric field frequency, the interaction between the energy levels decreases, and in the region of optical frequencies the energy levels become nearly isolated.

Then, it seems interesting to study the dependence of the transition probabilities on the electric field frequency and strength. To test the theoretical method accuracy, compare the calculated transition probabilities with the accurate data[17] in the absence of the electric field. The results of this comparison are presented in Table 1.

Table 1. Probabilities of the transitions $1snlLSJ - 1sn'l'L'SJ$ in the absence of electric field for the He atom, in $10^6 \cdot s^{-1}$

Transition $i \to k$	λ, nm	A_{ik}	
		This work	Ref. 17
7^3P_0–2^3S_1	276.462	1.17	1.11
8^3P_0–2^3S_1	272.400	0.79	0.78
9^3P_0–2^3S_1	269.692	0.57	0.55
5^1P_1–2^1S_0	361.467	3.80	3.76
6^1P_1–2^1S_0	344.858	2.27	2.39
7^3D_1–2^3P_0	370.620	2.33	2.47
9^1P_1–2^1S_0	325.921	0.69	0.65
6^3S_1–2^3P_0	386.873	0.27	0.26
8^3S_1–2^3P_0	365.316	0.11	0.12
8^3S_1–2^3P_1	365.303	0.32	0.36

As is seen, the calculated transition probabilities are in a good agreement with the accurate data, that indicates the reliability of the calculation technique. Table 2 summarizes the calculated probabilities of the transitions $nlJM$–$n'l'J'M'$ ($J, J' \le 1$) between magnetic sublevels for the He atom in the high-frequency discharge.

Table 2. Probabilities of the transitions $nlJM - n'l'J'M'$ in the high-frequency discharge (ω = 100 MHz), in $10^6 \cdot s^{-1}$

Transition $i \to k$	λ, nm (at $F = 0$)	A_{ik} (at $F = 0$)	$M \to M'$	A_{ik}	
				F, kV/cm	
				0.2	1
8^3P_0–2^3S_1	272.400	0.79	0 → −1	0.79	0.30
			0 → +1	0.79	0.28
9^3P_0–2^3S_1	269.692	0.57	0 → −1	0.56	0.40
			0 → +1	0.56	0.10
7^1P_1–2^1S_0	335.552	1.45	−1 → 0	1.46	1.41
			+1 → 0	1.46	1.42
8^1P_1–2^1S_0	329.772	0.98	−1 → 0	0.99	0.88
			+1 → 0	0.99	0.90
9^1P_1–2^1S_0	325.921	0.70	−1 → 0	0.70	0.51
			+1 → 0	0.70	0.52
5^3D_1–2^3P_0	402.749	6.83	−1 → 0	6.80	4.25
			+1 → 0	6.72	4.27
6^3D_1–2^3P_0	382.084	3.79	−1 → 0	2.08	1.59
			+1 → 0	3.22	1.02
7^3D_1–2^3P_0	370.620	2.33	−1 → 0	1.70	1.0
			+1 → 0	0.65	0.5

It follows from Table 2 that for all considered transitions the transition probabilities decrease with the increasing electric field strength. In addition, for all transitions, even if they have identical probabilities at a low field strength (F = = 0.2 kV/cm), the probabilities of the population of magnetic sublevels of the same level begin to differ with the increasing field strength. The difference between these probabilities can be quite significant, in particular, for 7^3D_1–2^3P_0 the transition probabilities between the magnetic sublevels differ twofold.

The results of investigation of the dependence of the transition probabilities between magnetic sublevels on the electric field frequency are given in Table 3 and Fig. 2.

Table 3. Probabilities of the transitions $9^3P_0(M)$–$2^3S_1(M')$ as functions of the frequency of the electric field, in $10^6 \cdot s^{-1}$

ω, MHz	$M \to M'$	A_{ik}	
		F = 0.2 kV/cm	F = 1 kV/cm
100	0 → −1	0.190	0.402
	0 → +1	0.189	0.103
$241.813 \cdot 10^5$	0 → −1	0.187	0.182
	0 → +1	0.187	0.190
$283.005 \cdot 10^6$	0 → −1	0.187	0.187
	0 → +1	0.187	0.187

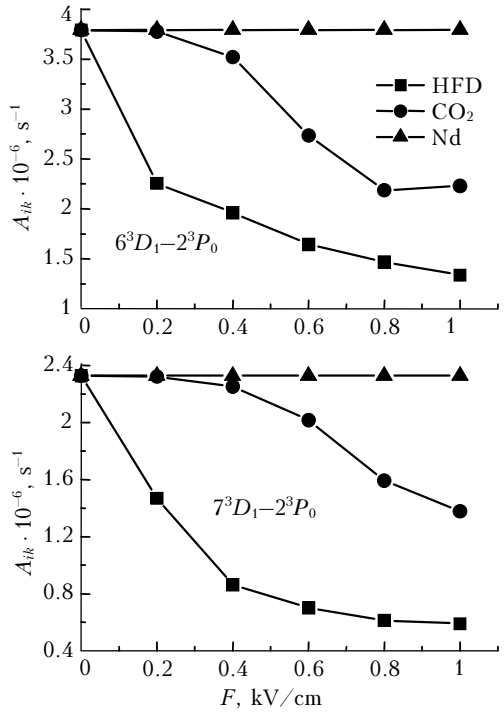

Fig. 2. Probabilities of transitions nd^3D_1–$2p^3P_0$ (n = 6,7) as functions of the electric field frequency: (HFD) high-frequency discharge, ω = 100 MHz; (CO_2) CO_2 laser, ω = = $241.813 \cdot 10^5$ MHz; (Nd) Nd laser, ω = $283.005 \cdot 10^6$ MHz.

It is seen from Table 3 that an increase of the electric field frequency leads to decrease in the anisotropy of transition probabilities between magnetic sublevels and, as follows from Table 3 and Fig. 2, to the lower sensitivity of the transition probability to the electric field strength. This loss in the sensitivity is caused by a weaker interaction between the energy levels, which is observed with the increasing electric field frequency. It should be noted that the transition probabilities have the polynomial dependence on the electric field strength, namely, $A_{ik} = a + bF + cF^2 + dF^3$ for the high-frequency discharge and $A_{ik} = a + bF + cF^2 + dF^3 + eF^4 + fF^5$ for the CO_2 laser. At the electric field frequency, corresponding to the Nd laser radiation, probabilities of transitions nd^3D_1–$2p^3P_0$ are insensitive to variations of the electric field strength in the considered range. Similar polynomial dependences are also observed for probabilities of other transitions.

Consider now the probabilities of transitions JM–$J'M'$ between magnetic sublevels at $J, J' \le 2$. The calculated probabilities of the transitions are shown in Fig. 3, from which it is seen that, as in the case of transitions between the levels with $J, J' \le 1$, the increase in the electric field strength leads to decrease of the probabilities. In addition, at all considered frequencies of the electric field, the strong anisotropy is observed in the transition probabilities: the larger the magnetic quantum number M for $M \to M'$, the higher the probability of the transition to the sublevel M. The maximal transition probability corresponds to max|M|. The most pronounced anisotropy of probabilities is observed in a high-frequency discharge as a consequence of the strongest (as compared to other electric field frequencies) interaction between the energy levels.

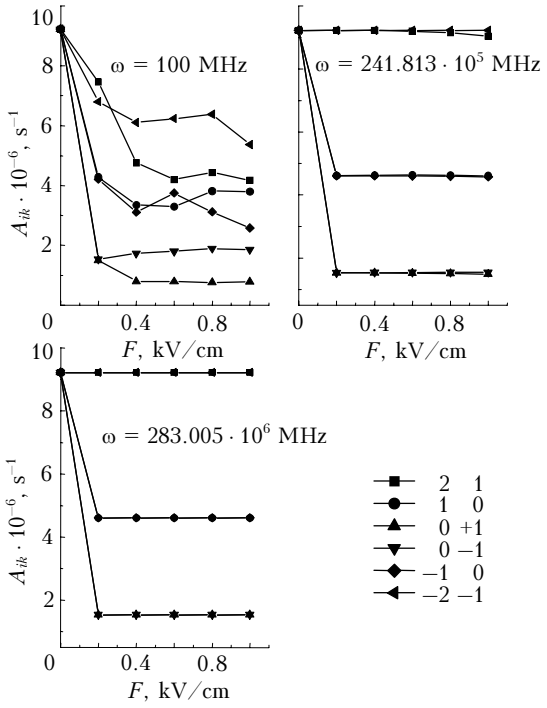

Fig. 3. Probabilities of the transition 5^3D_2–2^3P_1 as functions of the electric field frequency.

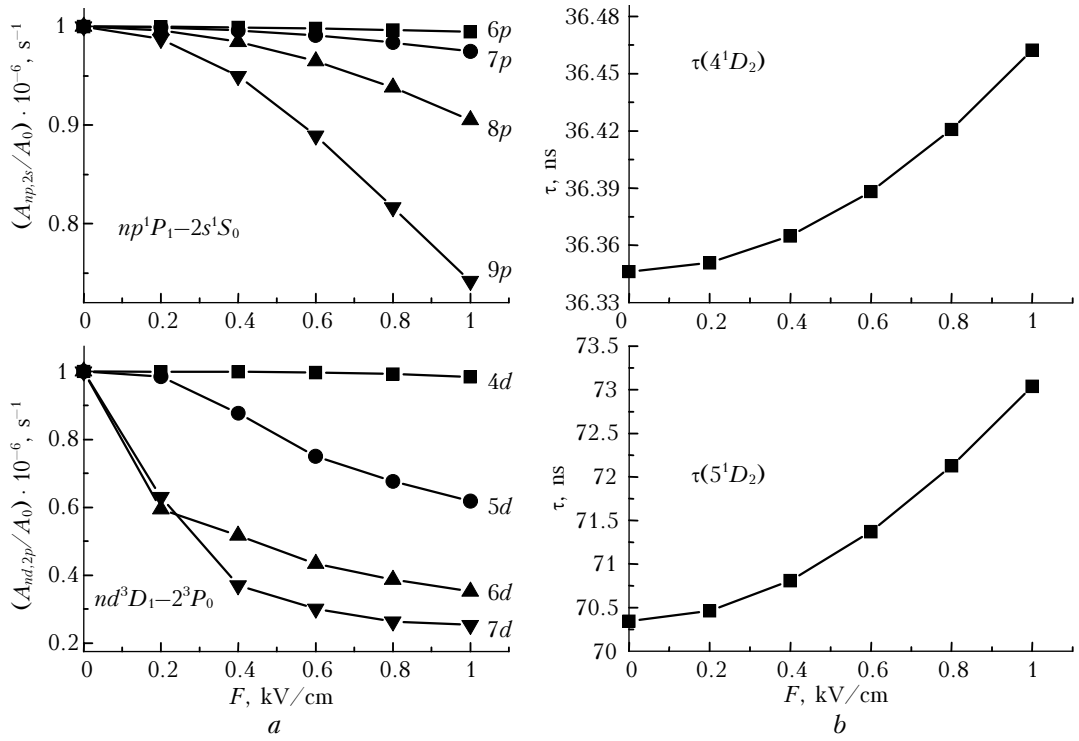

Fig. 4. Behavior of the transition probabilities and lifetimes of energy levels for the He atom in the high-frequency discharge: probabilities of the transitions nlJ–$n'l'J$ (*a*); lifetime of the nd^1D_2 states (n = 4, 5) (*b*).

As to the probabilities of the transitions J–J', calculated by Eq. (17), and the lifetimes of the He atom states in the external electric field, calculated by Eq. (18), they, obviously, also have the polynomial dependence on the electric field strength. For the considered range of the electric field strength, the following polynomial dependences were obtained: cubic polynomial for the high-frequency discharge, fifth-order polynomial for the CO_2 laser, and no dependence on the field strength for the Nd laser. As an illustration, Fig. 4 depicts the dependences of the probabilities of the transitions J–J' (Fig. 4*a*) and the lifetimes of the states nd^1D_2 (Fig. 4*b*) for the He atom in the high-frequency discharge. In Fig. 4*a*, $A_{JJ'}/A_0$ is the ratio of the transition probabilities $A_{JJ'}$ for the atom in the electric field to the probability A_0 of the corresponding transition in the absence of an electric field.

The reliability of the results obtained can be assessed by comparing the lifetimes calculated in the absence of the electric field with other theoretical results and with the experiment. The experimental data for lifetimes of the considered levels in the electric field are absent. The results of the comparison are given in Table 4.

Table 4. Lifetimes of the states nd^1D_2 of the He atom in the absence of electric field

τ, ns	This work	Theory [Ref. 18]	Experiment [Ref. 18]		
$\tau(4^1D_2)$	36.34	37	35±4	35±4	38±2
$\tau(5^1D_2)$	70.35	71.5	72±4	72±5	71.9±18

It is seen that the lifetimes calculated in this work are in a good agreement with the results obtained by other theoretical methods and with the experimental data. As can be seen from Fig. 4, for the high-frequency discharge both the transition probabilities and the lifetimes depend on the electric field strength as a cubic polynomial. The higher is the state, from which the transition occurs, the more sensitive is the transition probability and, consequently, the lifetime of the atomic state to variations in the field strength.

Conclusions

General analysis of calculations performed in this work reveals some regularities for the dynamic Stark effect, transition probabilities, and lifetimes for the He atom in the circularly polarized field of arbitrary strength and frequency.

First, the increase in the frequency of the electric field leads to a smaller split and shift of spectral lines.

Second, probabilities of the allowed transitions in most cases decrease with the increasing electric field strength. In addition, the forbidden lines appear due to interaction between Stark levels. The probability of the forbidden transitions increases with the increasing field strength. As the frequency of the electric field increases, the sensitivity of the transition probabilities to variations of the electric field strength decreases.

Third, the anisotropy of probabilities has been found for a significant number of transitions. In most

cases, it is connected with interaction of atomic energy levels in the electric field.

Finally, the polynomial dependence of the transition probabilities on the electric field strength has been revealed for all considered transitions.

The theoretical results obtained in this work can be used for explanation of physical processes occurring in the plasma generated in the circularly polarized alternating electric field. These results can also be useful for the modeling of new sources of light and excitation.

Acknowledgments

This work is supported, in part, by INTAS Grant No. 01–0200.

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