

PACS numbers: 32.80.–t; 42.50.Md
DOI: 10.1070/QE2005v035n06ABEH003422

Coherent phase control of excitation of atoms by bichromatic laser radiation in an electric field

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***Abstract*. A new method for coherent phase control of excitation of atoms in a discrete spectrum under the action of bichromatic laser radiation with the frequency ratio 1 : 2 is analysed. An important feature of this control method is the presence of a electrostatic field, which removes the parity selection rule for one of the control channels. It is shown that for the phase difference between the monochromatic radiation components, corresponding to the destructive interference between channels, there exists the electrostatic field strength at which the excited atomic transition is 'bleached'. It is proposed to use luminescence at the adjacent atomic transition for detecting the phase dependence of optical excitation.**

Keywords: *bichromatic radiation, coherent control, quantum interference*.

1. Introduction

The phase control of photoinduced processes in multicolour laser fields has been extensively studied in the last decade [1 – 4]. Interest in this problem is explained by the fact that the phase control of the rate of photostimulated effects is performed not by varying the field amplitude, as usual, but by changing the relative phase of monochromatic components of radiation.

The phase control is based on the quantum interference of the amplitudes of transitions corresponding to the different channels of the process, which are induced by the monochromatic components of radiation. A change in the relative phase of these components causes a change in the interference term in the total amplitude of a signal, resulting in the phase modulation of the process probability.

The idea of such a coherent control was first proposed by Russian scientists [5], who analysed the possibility of 'bleaching' of a medium upon a multiphoton resonance. Note that interest in this field was initiated by studies of the third harmonic generation in the case of a double-frequency resonance [6]. A related coherent photogalvanic effect was investigated in [7] as a mechanism of photoinduced second harmonic generation in optical fibres.

One of the first phase-control experiments in a bichromatic laser field was performed by Russian researchers [8], who observed the polar asymmetry of a photocurrent from a solid surface excited by fundamental and second-harmonic laser radiation. The same authors studied the coherent control of a photocurrent produced upon photoionisation of preliminary excited sodium atoms by bichromatic radiation with the frequency ratio 1 : 2 [9].

Note that in [8, 9] a bound – free electronic transition induced by laser radiation was observed. The interference between the quantum amplitudes of bound – bound transitions in a bichromatic laser field was first observed in [10] for the $6s\,^1S_0 \rightarrow 6p\,^1P_1$ transition in a mercury atom excited by three photons at the fundamental frequency and one photon of the third harmonic. The excitation probability of a mercury atom was modulated by varying the phase difference between the monochromatic components of radiation. The phase modulation of the excitation probability was detected by a signal of two-photon ionisation of the $6p\,^1P_1$ state. As a result, the possibility of coherent phase control of excitation of an atom in a discrete spectrum was convincingly demonstrated.

The phase control of atomic effects in bichromatic laser fields was later investigated in a number of theoretical and experimental papers. These studies (till the year 2000) are reviewed in [1]. Apart from photoionisation of atoms and photodecomposition of negative ions, the generation of harmonics, elastic and inelastic scattering of electrons, X-ray scattering, etc. were investigated. For example, the phase-dependent autoionisation of calcium atoms induced by the fundamental and third-harmonic radiation was observed in [11]. This opens up the possibility of the coherent control of the spectral shape of an autoionisation resonance line, which was earlier assumed completely determined by the structure of an atom.

The phase control was studied both in relatively weak (perturbation regime) [8 – 10] and strong fields (for example, above-threshold photoionisation of atoms [12]), when the perturbation theory cannot be applied. The papers on the coherent control on molecular dynamics by using bichromatic fields are presented in review [2]. Note in this connection the experiment [13] devoted to the phase control of the photodissociation of the HD^+ molecular ion induced by the 1053-nm fundamental radiation and its second harmonic.

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Received 15 December 2004; revision received 24 February 2005
Kvantovaya Elektronika **35** (6) 541 – 546 (2005)
Translated by M.N. Sapozhnikov

The coherent phase control of the internal photoeffect was also performed. In [14], a photocurrent in a GaAs bulk sample was measured, which was produced due to the quantum interference of one- and two-photon absorptions at the wavelengths 0.775 and 1.55 μm in the absence of an external potential difference. The monochromatic field does not produce a photocurrent in a centrally symmetric semiconductor due to the parity selection rule if an external electric voltage is not applied to the sample. In a bichromatic field, the photocurrent can be generated under these conditions, and its magnitude and direction can be controlled by varying the phase difference between the monochromatic components of radiation [14].

The studies devoted to the coherent control have been further developed in the last five years. Apart from processes in laser fields, the excitation and ionisation of Rydberg atoms induced by bichromatic gigahertz fields were studied theoretically and experimentally. Thus, the authors of paper [15] performed the phase control of the ionisation of an atomic hydrogen preliminary excited to the $n = 51$ Rydberg state under the action of a linearly polarised bichromatic field consisting of monochromatic components with frequencies 6 and 18 GHz. The phase dependence of the probability of excitation of a helium atom from the $28\,^3S$ Rydberg state by a bichromatic field of the same type as in [15] was studied experimentally in [16].

Let us emphasise that the phase control was performed in the above-mentioned studies by using laser pulses with duration greatly exceeding the optical cycle duration, when the absolute phase of the bichromatic field is insignificant and the probability of processes is determined by the phase difference of the monochromatic components. These studies are related to the analysis of optical processes induced by ultrashort laser pulses at one central frequency with duration of the order of a light wave cycle. The phase of the carrier frequency in such pulses is changed in a controllable way with respect to the pulse envelope, resulting in the modulation of the optical signal. In [17], the modulation of a photocurrent from the gold surface excited by a 750-nm, 4-fs laser pulse was studied as a function of the absolute phase of radiation.

The qualitative relation between effects existing in a bichromatic laser field and induced by ultrashort laser pulses with a variable absolute phase was obtained in [18], where the photoionisation of xenon atoms in the non-perturbation regime induced by a bichromatic field of a special type was calculated. The ionising field [18] consisted of the fundamental radiation and its second harmonic with a controllable relative phase, the field amplitude at the fundamental frequency being twice that of the second harmonic. The resulting field of such radiation represents an infinite train of pulses with a duration equal to the optical cycle. A change in the relative phase of the monochromatic components in the bichromatic field corresponds to a change in the absolute phase of the carrier frequency of these pulses, as in the experiment with femtosecond pulses [17].

Despite an increasing variety of the types of phase control of optical processes being studied at present, most of the investigations have been performed so far for the bound–free and free–free electronic transitions in a bichromatic field [1]. In this case, the parity selection rules do not prevent the phase modulation of the signal. Another situation appears upon excitation of atomic electrons in a discrete spectrum, when the selection rules impose certain restrictions on the frequencies of monochromatic components of radiation. The necessary condition is the same parity of the number of photons in each of the channels of the process. In [10], each of the two photoexcitation channels of a mercury atom contained one or three photons. In [19], the phase control of photoionisation of krypton and xenon atoms was performed through a resonance state, which was excited via two channels by four photons at the fundamental frequency and a photon at the fundamental frequency plus a photon at the third-harmonic frequency.

In this paper, a new type of phase control is proposed for excitation of atoms in a discrete spectrum by a bichromatic field consisting of the fundamental radiation and its second harmonic. To remove the parity selection rule, the atom being excited is subjected to the action of an external electrostatic field. In this case, not only phase control can be performed with the help of a bichromatic field of the specified type but also an additional control parameter is obtained – the constant field strength, which can be of interest for possible applications.

2. Calculation of the process probability

Consider the coherent phase control of excitation of a valence electron of an alkali metal atom in a discrete spectrum (Fig. 1). In this case, the process probability can be calculated most simply and the corresponding laser frequencies lie in the visible and near-UV regions.

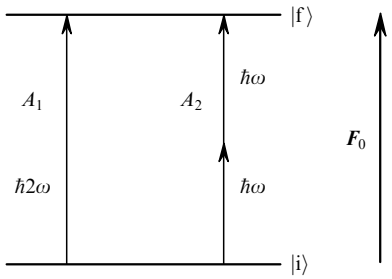

Figure 1. Two channels of the electronic transition in an atom from the $|i\rangle$ state to the $|f\rangle$ state with the quantum amplitudes A_1 (one-photon absorption at the second harmonic frequency) and A_2 (two-photon absorption at the fundamental frequency) caused by bichromatic radiation in the presence of a electrostatic field $\boldsymbol{F}_0$.

Let an alkali metal atom be subjected to the action of a bichromatic field representing a coherent superposition of the fundamental radiation at the frequency ω and its second harmonic at the frequency 2ω with a controllable difference of phases φ_1 and φ_2. The electric field strength vector of this radiation is

$$\boldsymbol{F}(t) = \boldsymbol{F}_1 \cos(\omega t + \varphi_1) + \boldsymbol{F}_2 \cos(2\omega t + \varphi_2). \quad (1)$$

We assume that the atom is located in a constant homogeneous electric field of strength $\boldsymbol{F}_0$, which mixes the electronic states of the atom with different parities.

We also assume here that the electric field strengths $\boldsymbol{F}_{0,1,2}$ are many orders of magnitude lower than the atomic electric field strength $F_a = 5.14 \times 10^9$ V cm^{-1}, so that the process

can be calculated by using the perturbation theory (perturbation regime). Then, we can neglect, with a reasonable accuracy, the 'dressing' of the atom by the electromagnetic field because this effect introduces a small correction to the process amplitude calculated in the first non-vanishing order of the perturbation theory. The application of the perturbation approach in this case is alleviated by the absence of one-photon resonance at which atomic states can be strongly mixed even in weak fields. The validity of this method is also confirmed by good agreement between experiments and the theory in numerous studies employing relatively weak laser fields (see, for example, [8–11, 19], in which the perturbation theory was used for calculations by neglecting the 'dressing' of atoms.

Let us calculate the excitation probability of an atom by the bichromatic field (1) per unit time taking into account the perturbation of its electronic states by a electrostatic field. We assume that the initial and final states of the atom have the zero orbital momentum (s states). The wave functions of the initial and final states of the atom perturbed by the constant homogeneous electric field of strength $\boldsymbol{F}_0$ in the first order of the perturbation theory with respect to the electric dipole interaction operator $\hat{V} = -\boldsymbol{d}\boldsymbol{F}_0$ are described by the expression

$$\Psi_{ns} = \Psi_{ns}^{(0)} + \sum_j \frac{\langle jp|\boldsymbol{d}\boldsymbol{F}_0|ns\rangle}{E_j - E_{ns}}\, \Psi_{jp}^{(0)}, \tag{2}$$

where $\Psi_j^{(0)}$ are the unperturbed wave functions of the atom; E_j are the atomic energy levels; and $\boldsymbol{d}$ is the dipole moment operator of the atom. Expression (2) takes into account that only transitions from the s states to the p states of the atom are allowed by the selection rules. The fine splitting of atomic energy levels is neglected.

The matrix element of the dipole moment operator calculated between the perturbed wave functions (2) of the initial $|n_\mathrm{i}\mathrm{s}\rangle$ and final $|n_\mathrm{f}\mathrm{s}\rangle$ states of the atom is

$$\langle n_\mathrm{f}\mathrm{s}|\boldsymbol{d}|n_\mathrm{i}\mathrm{s}\rangle = F_{0k} \times \sum_j \left\{ \frac{\langle n_\mathrm{f}\mathrm{s}|d_k|jp\rangle\langle jp|\boldsymbol{d}|n_\mathrm{i}\mathrm{s}\rangle}{E_j - E_{n_\mathrm{f}\mathrm{s}}} + \frac{\langle n_\mathrm{f}\mathrm{s}|\boldsymbol{d}|jp\rangle\langle jp|d_k|n_\mathrm{i}\mathrm{s}\rangle}{E_j - E_{n_\mathrm{i}\mathrm{s}}} \right\}. \tag{3}$$

Here, summation is performed over twice repeated subscripts.

The absorption amplitude A_1 for one second-harmonic photon at the $|n_\mathrm{i}\mathrm{s}\rangle \to |n_\mathrm{f}\mathrm{s}\rangle$ transition, taking into account the perturbation of the atom by the electrostatic field, has the form

$$A_1 = -\frac{\langle n_\mathrm{f}\mathrm{s}|\boldsymbol{d}|n_\mathrm{i}\mathrm{s}\rangle \boldsymbol{F}_2}{\hbar} = \frac{1}{2\hbar} F_{0k}F_{2l} \exp(-\mathrm{i}\varphi_2) \times \sum_j \left\{ \frac{\langle n_\mathrm{f}\mathrm{s}|d_k|jp\rangle\langle jp|d_l|n_\mathrm{i}\mathrm{s}\rangle}{E_j - E_{n_\mathrm{f}\mathrm{s}}} + \frac{\langle n_\mathrm{f}\mathrm{s}|d_l|jp\rangle\langle jp|d_k|n_\mathrm{i}\mathrm{s}\rangle}{E_j - E_{n_\mathrm{i}\mathrm{s}}} \right\}. \tag{4}$$

Expression (4) takes into account that absorption is induced by the positive-frequency component of the second-harmonic electric field, whose complex amplitude is $\tilde{\boldsymbol{F}}_2 = (\boldsymbol{F}_2/2)\exp(-\mathrm{i}\varphi_2)$. The sum over the intermediate states in the right-hand side of equality (4) can be expressed in terms of the scattering tensor of the electromagnetic field upon excitation of the atom from the initial $|n_\mathrm{i}\mathrm{s}\rangle$ state to the final $|n_\mathrm{f}\mathrm{s}\rangle$ state

$$c_\mathrm{fi}^{lk}(\omega, \omega') = \frac{1}{\hbar}\sum_j \left\{ \frac{d_{fj}^l d_{ji}^k}{\omega_{ji} - \omega - \mathrm{i}0} + \frac{d_{fj}^k d_{ji}^l}{\omega_{ji} + \omega' - \mathrm{i}0} \right\}. \tag{5}$$

By comparing the expression for the sum in (4) and the definition of the scattering tensor (5), we obtain

$$\sum_j \left\{ \frac{\langle n_\mathrm{f}\mathrm{s}|d_k|jp\rangle\langle jp|d_l|n_\mathrm{i}\mathrm{s}\rangle}{E_j - E_{n_\mathrm{f}\mathrm{s}}} + \frac{\langle n_\mathrm{f}\mathrm{s}|d_l|jp\rangle\langle jp|d_k|n_\mathrm{i}\mathrm{s}\rangle}{E_j - E_{n_\mathrm{i}\mathrm{s}}} \right\} = c_\mathrm{fi}^{lk}(\omega_\mathrm{fi}, 0). \tag{6}$$

According to the law of conservation of energy, we have the equality

$$\hbar\omega_\mathrm{fi} = 2\hbar\omega. \tag{7}$$

As a result, the amplitude of one-photon absorption of radiation at the second-harmonic frequency is described by the expression

$$A_1 = \frac{1}{2\hbar}(\boldsymbol{F}_0\boldsymbol{F}_2)\exp(-\mathrm{i}\varphi_2)c_\mathrm{fi}(2\omega, 0). \tag{8}$$

Expression (8) takes into account that the scattering tensor for transitions between atomic states with the zero momentum transforms to the scalar: $c_\mathrm{fi}^{lk} = c_\mathrm{fi}\delta^{lk}$.

The expression for the amplitude of two-photon absorption of radiation at the fundamental frequency can be obtained similarly [20]. It has the form

$$A_2 = \frac{1}{4\hbar}\boldsymbol{F}_1^2 \exp(-2\mathrm{i}\varphi_1)c_\mathrm{fi}(\omega, -\omega). \tag{9}$$

The total excitation probability of an atom per unit time is proportional to the square of the modulus of the sum of amplitudes (8) and (9) and has the form

$$w_\mathrm{fi}^\mathrm{tot}(\varphi) = \frac{4\pi^3}{\hbar^2 c^2} G_\mathrm{fi}(\omega) I_1^2 \left| c_\mathrm{fi}(\omega, -\omega) + 2\cos(\alpha) r_{01}\sqrt{\eta}\, \mathrm{e}^{\mathrm{i}\varphi} c_\mathrm{fi}(2\omega, 0)\right|^2, \tag{10}$$

where $r_{01} = F_0/F_1$ is the ratio of the strengths of the electrostatic field and the electric field of radiation at the fundamental frequency; $I_1 = cF_1^2/8\pi$ is the radiation intensity at the fundamental frequency; $\eta = I_2/I_1$ is the intensity ratio for the monochromatic components of the field (1); $G_\mathrm{fi}(\omega)$ is the shape of the $|\mathrm{i}\rangle \to |\mathrm{f}\rangle$ transition line; α is the angle between the vectors $\boldsymbol{F}_0$ and $\boldsymbol{F}_2$; and $\varphi = 2\varphi_1 - \varphi_2$ is the controllable phase shift between the monochromatic components of radiation (2) ('bichromatic' phase).

The first term in the modulus in (10) corresponds to two-photon excitation of the atom at the fundamental frequency, while the second one corresponds to one-photon absorption of the second harmonic by the atom perturbed by the electrostatic field. Let us emphasise that the second term in the absence of a constant field is zero for the transitions between the s states of the atom.

Strictly speaking, expression (10) was obtained for continuous bichromatic radiation. A finite duration of laser field pulses can drastically change the phase dependence (10) only in the case of ultrashort (femtosecond) pulses of duration comparable with the light wave cycle. Then, as pointed out in Introduction, the process probability begins

to depend not only on the phase difference between the monochromatic components but also on the absolute phase of radiation. In the case of picosecond and nanosecond pulses, expression (10) is modified to take into account the time dependence of the radiation intensity, and the right-hand side of equality (10) is convoluted with the laser line shape.

In calculations by expression (10), it is convenient to express the scattering tensor of the atom in terms of the transition oscillator strengths,

$$c_{\mathrm{fi}}(\omega,\omega') = \frac{e^2}{2m}\sum_j \left(\frac{f_{j\mathrm{i}} f_{j\mathrm{f}}}{\omega_{j\mathrm{i}}\omega_{j\mathrm{f}}}\right)^{1/2} \times\left(\frac{1}{\omega_{j\mathrm{i}}-\omega-\mathrm{i}0}+\frac{1}{\omega_{j\mathrm{i}}+\omega'-\mathrm{i}0}\right), \tag{11}$$

because the oscillator strengths $f_{j\mathrm{i(f)}}$, as the eigenfrequencies $\omega_{j\mathrm{i}}$, are tabulated [21, 22].

Let us transform the right-hand side of equality (10) to the form

$$w_{\mathrm{fi}}^{\mathrm{tot}}(\varphi) = w_{\mathrm{fi}}^{\Sigma}\left[1+\frac{\zeta(\eta,r_{01})}{2}\cos\varphi\right], \tag{12}$$

where w_{fi}^{Σ} is the excitation probability of the atom neglecting the interference term and $\zeta(\eta,r_{01})$ is the depth of phase modulation. By using (10), we define the latter parameter by the expression

$$\zeta(\eta,r_{01}) = \frac{8\cos\alpha r_{01}\sqrt{\eta}\,c_{\mathrm{fi}}(\omega,-\omega)c_{\mathrm{fi}}(2\omega,0)}{c_{\mathrm{fi}}^2(\omega,-\omega)+4(\cos^2\alpha)r_{01}^2\eta c_{\mathrm{fi}}^2(2\omega,0)}. \tag{13}$$

For the optimum values of $r_{01}=F_0/F_1$ and $\eta=I_2/I_1$ obtained from the condition of the maximum phase modulation of the process ($\zeta=2$), we have

$$\left(|\cos\alpha|\sqrt{\eta}\,r_{01}\right)_{\mathrm{opt}} = \left|\frac{c_{\mathrm{fi}}(\omega,-\omega)}{2c_{\mathrm{fi}}(2\omega,0)}\right|, \tag{14}$$

where the frequency ω satisfies the law of conservation of energy (7).

3. Coherent control of photoexcitation

Let us use the expressions obtained in the previous section to analyse the possibility of the phase control of photoexcitation in a bichromatic laser field (1). Consider, for example, the 3s → 4s transition from the ground to the first excited state with the zero moment in a sodium atom. In this case, two-photon absorption at the fundamental frequency is allowed by the dipole selection rules. The parity selection rule for one-photon absorption of the second harmonic is removed by applying a electrostatic field. The numerical value of the right-hand side of equality (14) for excitation of the transition under study equal to 42.4 can be easily calculated from expression (11) using the oscillator strengths and transition frequencies presented in [21, 22]. In this case, the resonance wavelength λ_1 at the fundamental frequency for which the law of conservation of energy (7) is fulfilled is equal to 777 nm.

Note that the Doppler effect, which is inherent in experiments with alkali metal vapours, will not affect noticeably the phase control method considered here. Because the Doppler frequency shift is much smaller than characteristic atomic frequencies, the tensors of radiation scattering from an atom, determining the phase dependence of the excitation probability (10), do not virtually change. The modification of expression (10) will be reduced to the consideration of the Doppler shift in the $|\mathrm{i}\rangle \to |\mathrm{f}\rangle$ transition line shape by making the corresponding frequency replacement and to the integration of the obtained expression with the velocity distribution function of atoms, which will not change the phase dependence of the process.

Figure 2 shows the phase modulation depth of the excitation probability of the 3s → 4s transition in a sodium atom as a function of the electric field strength ratio $r_{01}=F_0/F_1$ calculated from (14) for different values of the parameter $\eta=I_2/I_1$ for the zero angle between vectors $\boldsymbol{F}_0$ and $\boldsymbol{F}_2$. One can see that the optimum value of the electric field strength ratio r_{01}, at which the phase modulation of excitation is maximal, decreases with increasing parameter η.

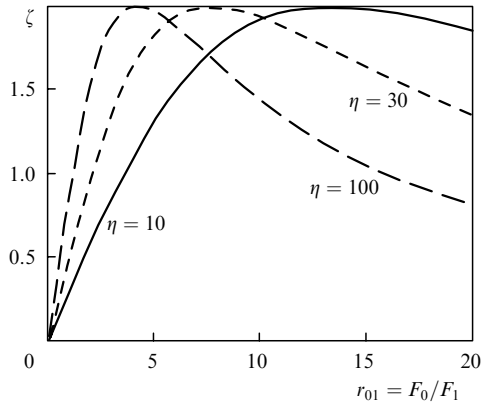

Figure 2. Parameter of phase modulation required to excite the 3s → 4s transition in a sodium atom by bichromatic radiation as a function of the strength ratio of a electrostatic field and the electric field of fundamental radiation for different values of $\eta=I_2/I_1$ and $\alpha=0$.

Depending on the value of the bichromatic phase $\varphi=2\varphi_1-\varphi_2$, interference between channels upon photoexcitation can be either constructive or destructive. Figure 3 shows the total probability of excitation of the 4s state in a sodium atom by the bichromatic field (1) as a function of a electrostatic field strength for different values of the bichromatic phase. The solid curve corresponds to the case $\varphi=\pi$, when interference between channels is destructive (the amplitudes of two-photon and one-photon processes are subtracted). In the case $\varphi=2\pi$, interference is constructive, and the amplitudes of channels are added. And finally, when $\varphi=\pi/2$, the interference term is absent.

It follows from Fig. 3 that in the case of destructive interference ($\varphi=\pi$), there exists the electric field strength at which the total excitation probability of the resonance transition per unit time vanishes. In other words, the bichromatic phase can be selected so that the electronic transition in the atom will be bleached for bichromatic radiation due to variations in the electrostatic field strength. If the bichromatic phase is not strictly equal to π (or 2π), but lies in the interval $\pi/2<\varphi<3\pi/2$, where the cosine is negative, the excitation probability will not be exactly zero.

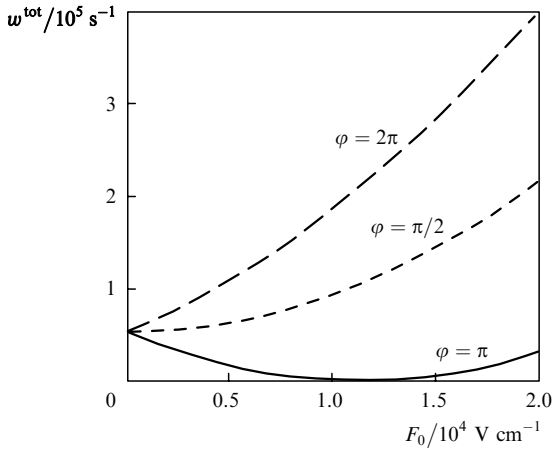

Figure 3. Total probability of excitation of the 3s → 4s transition in a sodium atom by bichromatic radiation per unit time as a function of a electrostatic field strength for different values of the bichromatic phase $\varphi = 2\varphi_1 - \varphi_2$ and $I_1 = 10^4$ W cm^{-2}, $\eta = 100$ and $\alpha = 0$.

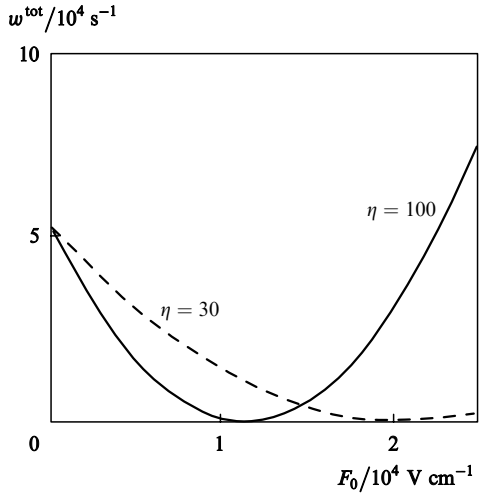

Figure 4. Dependences of the probability of excitation of a sodium atom at the 3s → 4s transition by a bichromatic field with the frequency ratio 1 : 2 on the electrostatic field strength for different $\eta = I_2/I_1$ and $I_1 = 10^4$ W cm^{-2}.

Instead, the curve $w^{tot}(F_0)$ will have a minimum, which is the deeper, the closer φ is to π.

The value of the 'bleaching' electrostatic field strength depends on the intensity ratio of the fundamental and second harmonic radiation. This dependence is shown in Fig. 4. One can see that the higher $\eta = I_2/I_1$, the lower the bleaching electrostatic field strength. This also follows from expression (14).

The phase control of excitation of an atom in a discrete spectrum by a bichromatic field can be observed by photoluminescence due to radiative transitions from an excited state of the atom to the lower energy levels. The corresponding scheme of luminescence upon excitation of the $|ns\rangle \rightarrow |n+1s\rangle$ transition in an alkali metal atom is shown in Fig. 5. The power of photoluminescence emitted from the excited $|f\rangle$ state of the atom in the unit volume into the solid angle $d\Omega$ is

$$\frac{dQ_{fj}}{dV d\Omega} = N_f A_{fj} \frac{\hbar\omega_{fj}}{4\pi}, \qquad (15)$$

where N_f is the concentration of atoms in the $|f\rangle$ state; A_{fj} is the $|f\rangle \rightarrow |j\rangle$ transition probability per unit time (the Einstein coefficient for spontaneous radiation); and $\hbar\omega_{fj}$ is the transition energy. A similar expression with the appropriate interchange of subscripts is valid for luminescence at the $|j\rangle \rightarrow |i\rangle$ transition. Because the upper level population is determined by the excitation probability $w^{tot}(\varphi)$ depending on the bichromatic phase $\varphi = 2\varphi_1 - \varphi_2$, the photoluminescence power can be used, according to (15), to observe the phase dependence of excitation of an atom in a discrete spectrum.

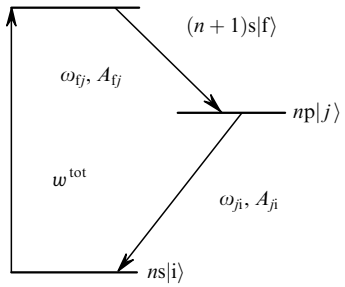

Figure 5. Scheme of excitation of photoluminescence of alkali atoms in the phase control method considered in the paper.

4. Conclusions

We have proposed a new method of phase-dependent excitation of an atom in a discrete spectrum by bichromatic excitation, which is based on the removal of the parity selection rule for dipole transitions by applying a electrostatic field. This allows one to perform the phase control of excitation of the atom due to the quantum interference of amplitudes of channels with the different parity of the number of photons involved in the process in different channels. In particular, it is possible to perform the phase modulation of excitation of an atom in a discrete spectrum by bichromatic radiation of the simplest type with the frequency ratio of monochromatic components equal to 1 : 2.

An interesting feature of this phase control method is the possibility of 'bleaching' of the excited transition by a proper tuning of the electric filed strength in the case of the destructive interference of the channels, when the bichromatic phase is equal to π.

In previous papers on the phase control of excitation of atoms, the effect was detected by the photoionisation of the excited state (see, for example, [10, 11]). In this paper, it is shown that the phase dependence of photoexcitation of atoms can be observed by luminescence emitted due to radiative transitions from the excited state to lower energy levels. The luminescent method for detecting the bichromatic phase can be promising for technical applications.

In conclusion, note that the phase control in a bichromatic field was performed in this paper by applying a electrostatic field. This goal can be also achieved by using radiation whose frequency is lower than the linewidth of the excited transition. In this case, the principal aspect of the problem remains the same (the removal of the parity selection rule), but new interesting effects can appear from the point of view of practical applications.

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