$t \approx 0.79T$. This suggests that the movement of the wave packet under the conditions considered is weakly relativistic, which justifies the use of the Schrödinger equation. On the other hand, quantum diffusion plays an important role in the process of generation of high-energy photons in recollisions. The role of diffusion proves to be twofold: although it reduces the electron density in the interaction region, at the same time it partially compensates for the action of the factors (the magnetic field of the wave, the ellipticity of polarization) that cause deviations of electron trajectories. As a result, although before a collision the trajectory of the electron wave packet can deviate from linear by a distance greatly exceeding the atomic size, a complete break of the generation of high-frequency photons does not occur.

6. Conclusions

In this paper, we have suggested and theoretically studied the method of generation of attosecond pulses based on the transformation of powerful femtosecond laser pulses on rapidly ionized atoms. The problems of generation of attosecond radiation are directly included in the general problem of nonlinear optical interactions of the pulses containing a small number of field periods. The theoretical and experimental studies of this problem are still in the initial stages. The model of the interaction of the optical field with a gas of quantum particles suggested and studied in this paper is sufficiently versatile to be applied to studies of various nonlinear optical phenomena in variety of media.

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Stabilization of atoms in a strong laser field

M V Fedorov

1. Introduction

The phenomenon of stabilization of atoms in a strong laser field was first predicted and described theoretically [1, 2] and only recently confirmed experimentally [3, 4]. The idea of stabilization of atoms does not correspond to the simple intuitive notion of an increase in the ionization rate of atoms with increasing intensity of the laser field. By definition, stabilization means that, depending on the intensity I or the strength ε of the field, beginning from some critical values of I_c or ε_c, the ionization rate Γ or the total ionization probability w_i during the pulse ceases to increase (i.e., is saturated at the level $w_i \ll 1$) or even becomes a decreasing function of I or ε. In this case, the ionization time t_i of an atom or the residual probability of its 'non-ionization' $w_{res} = 1 - w_i$ also either becomes constant or starts to increase with I or ε (Fig. 1). The choice between these types of behavior of the functions $w_{res}(\varepsilon)$, $\Gamma(\varepsilon)$, and $w_i(\varepsilon)$ in a strong field and also the determination of the critical fields I_c and ε_c is the object of contemporary studies of the physics of stabilization of atoms.

There are several models of stabilization of atoms in a strong field, which differ in both the physical nature and conditions of their realization. Below, we will briefly describe two of these models: the high-frequency stabilization (or the Kramers – Henneberger stabilization [1]) and the interference stabilization of atoms [2]. However, before proceeding to an analysis of them, we briefly consider the experimental data available.

2. Experiments

Actually, several attempts have been made to observe the stabilization of atoms induced by a strong field. However, it seems that only two of them [3, 4] can be considered fully successful. In both cases, the electron yield (proportional to the ionization probability w_i per pulse) was measured as a function of the pulse duration for a fixed fluence, i.e., for the product of the peak intensity I_{max} of the field with the pulse duration τ equal to $F = I_{max}\tau = \text{const}$; or as a function of F for different but fixed laser pulse durations τ. The objects of the studies were excited barium [3] and neon [4] atoms. Despite substantial differences in the setting up and interpretation of experiments [3, 4], their results are quite similar as a whole; therefore, we will describe only experiment [3].

The Ba atoms were prepared in the 27D state using two-stage excitation and were then ionized by a short powerful laser pulse. Figure 2a shows the electron yield measured as a function of the pulse duration $N_e(\tau)$ for $F = \text{const} = 7.8$ J cm^{-2}. A plateau, where $N_e(\tau) \approx \text{const}$, corresponds to the electron yield calculated in the first order of the perturbation theory using Fermi's 'golden' rule (FGR), according to which $N_e(\tau) \propto I_{max}\tau \equiv F = \text{const}$. The region

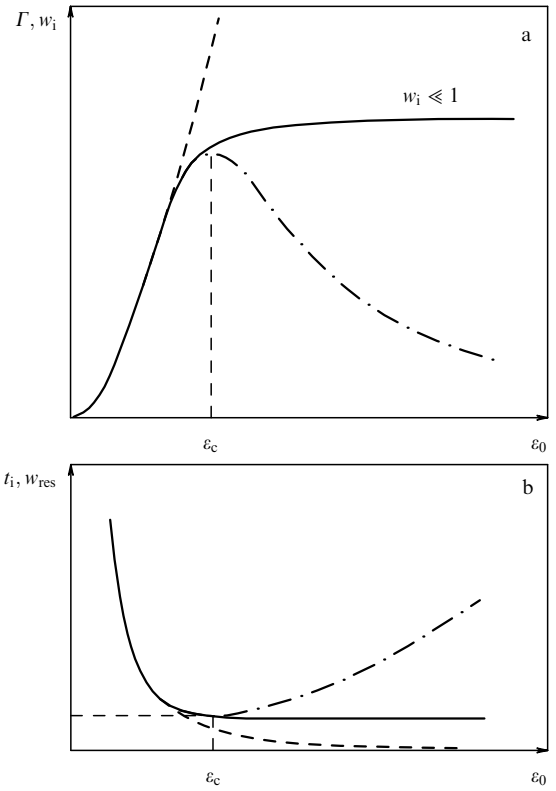

Figure 1. Dependences of (a) the probability w_i and the ionization rate Γ of an atom and (b) the residual probability w_{res} and the ionization time t_i on the peak electric strength of the laser pulse. The dashed line corresponds to the perturbation theory; the dashed-dotted and solid lines correspond to different regimes of stabilization of an atom in a strong field $\varepsilon_0 > \varepsilon_c$.

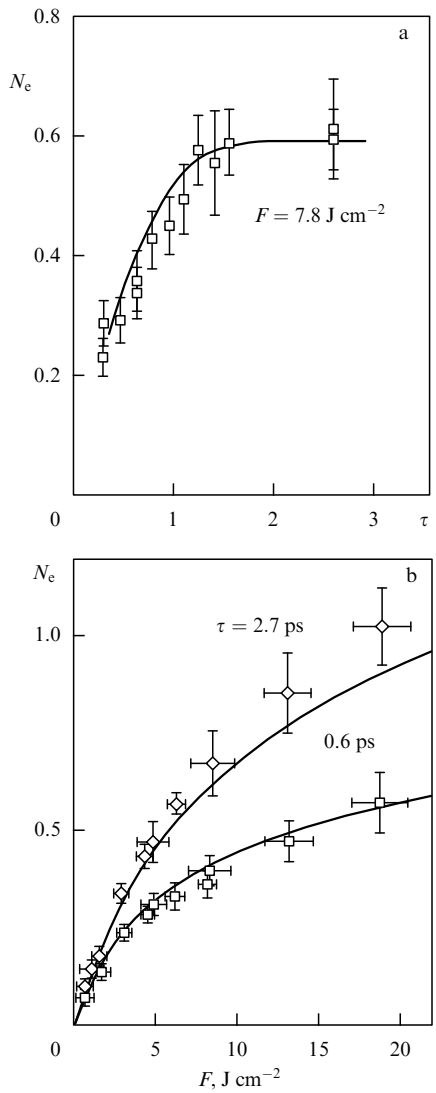

Figure 2. The electron yield measured as a function of the pulse duration τ at the constant illumination intensity F and as a function of the illumination intensity for the specified pulse duration τ [3].

of small τ for $F = \text{const}$ corresponds to high field intensities. The considerable deviation from FGR observed in this case [the decrease in $N_e(\tau)$] clearly confirms the effect of stabilization of the atom in a strong field. Figure 2b shows the experimental dependence of the electron yield on the fluence F for two fixed pulse durations $\tau = 2.7$ and 0.6 ps. Although the curve $N_e(F)$ related to the shorter pulse corresponds to the higher intensity I, it is located below the curve related to ionization by a longer pulse corresponding to the lower intensity I; this also confirms the presence of stabilization of atoms in a strong field. In addition, the form of curves $w_i(F)$ suggests that stabilization occurs due to saturation of the ionization probability $w_i(F) \approx \text{const}$ at the level $w_i \ll 1$.

3. The Kramers – Henneberger stabilization

We begin the description of various models and mechanisms of stabilization with the so-called Kramers – Henneberger model [1, 5 – 8]. It is known [9] that the Hamiltonian of an atom in an external light field of strength $\varepsilon = \varepsilon_0 \cos(\omega t)$ can be transformed in the dipole approximation to the form

$$H = \frac{\mathbf{p}^2}{2m} + U\big(\mathbf{r} - \boldsymbol{\alpha}_0 \cos(\omega t)\big)\,, \tag{1}$$

where $U(\mathbf{r})$ is the atomic potential in the absence of the field and $\boldsymbol{\alpha}_0 = \boldsymbol{\varepsilon}_0/\omega^2$ is the amplitude of oscillations of a free electron in the field. The transformation of the Hamiltonian to the form (1) is physically equivalent to transition to the oscillating coordinate system in which the electron is at rest, whereas the atomic nucleus and the corresponding atomic potential oscillate.

The potential in Eqn (1), which harmonically depends on time, can be expanded in a Fourier series

$$U\big(\mathbf{r} - \boldsymbol{\alpha}_0 \cos(\omega t)\big) = \sum_n U_n(\mathbf{r}; \boldsymbol{\alpha}_0) \exp(\mathrm{i}n\omega t)\,. \tag{2}$$

The Kramers – Henneberger approximation is based on the assumption that the term with $n = 0$ in expansion (2) plays the dominant role, whereas all harmonics (terms with $n \neq 0$) can be treated as a small perturbation. In this approximation, an atom in a strong laser field is characterized by the Kramers – Henneberger Hamiltonian H_{KH} with the stationary potential

$$U_{\mathrm{KH}}(\mathbf{r}) \equiv U_0(\mathbf{r}; \boldsymbol{\alpha}_0) = \frac{\omega}{2\pi} \int_0^{2\pi/\omega} \mathrm{d}t \exp(-\mathrm{i}n\omega t) \times U\big(\mathbf{r} - \boldsymbol{\alpha}_0 \cos(\omega t)\big)\,. \tag{3}$$

The main criteria for the validity of the Kramers–Henneberger approximation are

$$\alpha_0 \gg \Delta r_0 \quad \text{and} \quad \varepsilon_0 \gg \varepsilon_{\text{BSI}}\,, \tag{4}$$

where Δr_0 is the size of the localization region of the wave function of the unperturbed atomic electron and ε_{BSI} is the characteristic strength of the field in the model of the barrier suppression ionization (BSI) of an atom. This model and the determination of ε_{BSI} are illustrated in Fig. 3, where the total potential energy $U(x) = -e^2/|x| - \varepsilon x$ of a one-dimensional atom in an external static field is shown. The value of ε_{BSI} is determined from the condition that the maximum $U_{\max}$ of the curve $U(x)$ drops to the ground-state energy level E_0 of the atom. The BSI model assumes that in a strong alternating field, when its strength achieves the value ε_{BSI}, the atomic wave packet also penetrates over the barrier into the region far removed from the nucleus and behaves as the wave packet of a free particle in the field of the light wave. In the hydrogen atom, ε_{BSI} is equal to 1/16 of the atomic field strength, which corresponds to an intensity $I \sim 10^{14}$ W cm^{-2}. When the conditions (4) of applicability of the Kramers–Henneberger approximation are satisfied, the role of the harmonics with $n \neq 0$ in expansion (2) decreases with increasing field in a certain range of intensities, which suggests that the atom is stabilized. The condition $\varepsilon > \varepsilon_{\text{BSI}}$ determines the intensity I_c of the critical field required for the Kramers–Henneberger stabilization of an atom.

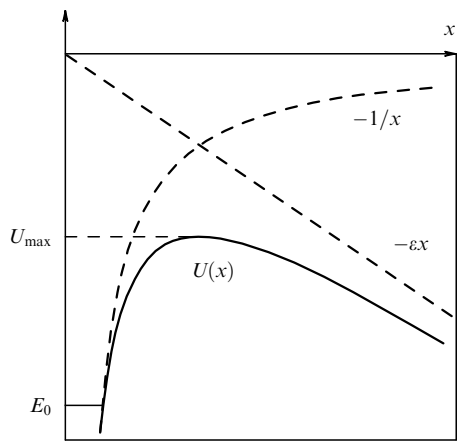

Figure 3. Potential energy $U(x)$ of an electron in the Coulomb and permanent electric fields (the one-dimensional model of the hydrogen atom); the field strength ε_{BSI} corresponding to the barrier suppression is determined by the condition $U_{\max} = E_0$.

4. Interference stabilization of Rydberg atoms

The interference stabilization of atoms [2, 9–16] is qualitatively different from the Kramers–Henneberger stabilization both in its physics and the conditions required for its realization. The main idea of interference stabilization is illustrated in Fig. 4. If an atom is initially prepared in a highly excited (Rydberg) state with energy E_n ($n \gg 1$), then the rate Γ of photoionization by the field with frequency $\omega > |E_n|$ is determined by the formula of the first-order perturbation theory (FGR) and linearly depends on the field intensity. The magnitude of Γ also determines the ionization

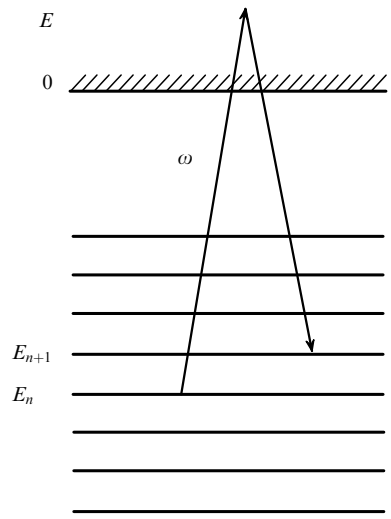

Figure 4. Raman transitions of the Λ type between Rydberg levels of an atom responsible for the appearance of the interference stabilization.

broadening of Rydberg levels (Fig. 4). The criterion for applicability of the perturbation theory and FGR is the smallness of the width Γ as compared to the distance between neighboring Rydberg levels, $\Gamma \ll E_{n+1} - E_n \approx 1/n^3$. In stronger fields (for $\Gamma \geqslant 1/n^3$), the perturbation theory becomes inapplicable, and along with direct transitions to a continuum, Raman transitions accompanied by the repopulation of neighboring Rydberg levels (the Λ transitions shown schematically in Fig. 4) become efficient. The coherent populations of Rydberg levels, which appear in this case, prove to be phased in such a way that the subsequent transitions from these levels to the continuum interfere and partially cancel each other, resulting in a decrease in the ionization rate of an atom and its stabilization. When quasiclassical expressions for matrix elements are used [16], the stabilization criterion $\Gamma \geqslant 1/n^3$ is reduced to the form $\varepsilon_0 \geqslant \omega^{5/3}$. In the optical frequency range ($\omega \sim 10^{-1}$), this condition corresponds to an intensity $I \geqslant 10^{13}$ W cm^{-2}, in accordance with experimental conditions [3]. Below, we discuss the main methods of theoretical description of the interference stabilization and their results.

Model of equations for the probability amplitudes of finding an atom in Rydberg states. The exact wave function $\Psi(\mathbf{r}, t)$ of an atom can be expanded in the wave functions of the continuous $[\psi_E(\mathbf{r})]$ and discrete $[\psi_n(\mathbf{r})]$ spectra with expansion coefficients $C_E(t)$ and $C_n(t)$. With the help of an approximation known as the adiabatic elimination of the continuum (see pp. 366–368 in Ref. [9]), the system of equations for the probability amplitudes $C_n(t)$ and $C_E(t)$ can be reduced to the system of equations only for functions $C_n(t)$. The simplest model that can be used for the description of interference stabilization is a model of two closely spaced levels E_n and E_{n+1} interacting with a continuum. Of course, such an approximation for the Rydberg atom serves only as an illustration, but it is very informative. As a result, a system of two equations for functions $C_\alpha(t)$ ($\alpha = n, n+1$) appears, in which the coupling with the continuum is determined by the tensor of ionization widths $\Gamma_{\alpha,\beta}$:

$$\mathrm{i}\dot{C}_\alpha - E_\alpha C_\alpha = -\frac{\mathrm{i}}{2}\sum_\beta \Gamma_{\alpha,\beta} C_\beta\,. \tag{5}$$

In the approximation of equal ionization widths ($\Gamma_{\alpha,\beta} \equiv \Gamma$), the quasi-energy solutions $C_\alpha(t) \propto \exp(-\mathrm{i}\gamma t)$ of Eqns (5) are very simple and, in particular, the complex quasi-energies $\gamma_\pm$ of this system have the form

$$\gamma_\pm = \frac{1}{2}\left[E_n + E_{n+1} - \mathrm{i}\Gamma \pm \sqrt{(E_{n+1} - E_n)^2 - \Gamma^2}\right]. \qquad (6)$$

Figure 5 shows the broadened quasi-energy levels (bands) of a two-level system in a strong ionizing field. The centers of mass of the bands and their widths are determined by the real (γ_+) and imaginary (γ_-) parts of the quasi-energies, respectively. Both quasi-energy bands broaden depending on Γ (i.e., on the field intensity I) as long as $\Gamma < \Delta \equiv E_{n+1} - E_n$. Then, for $\Gamma > \Delta$, one of the bands continues to broaden, whereas another band narrows. The appearance of the narrow quasi-energy level (band) γ_- in a strong field suggests that the system can be stabilized: the lifetime of the long-lived component of the population is determined in this case by the value of $1/|\gamma''_-|$. The centers of mass of both bands in a strong field (for $\Gamma > \Delta$) are localized midway between the energies of the unperturbed levels: $\gamma'_\pm = (E_n + E_{n+1})/2$.

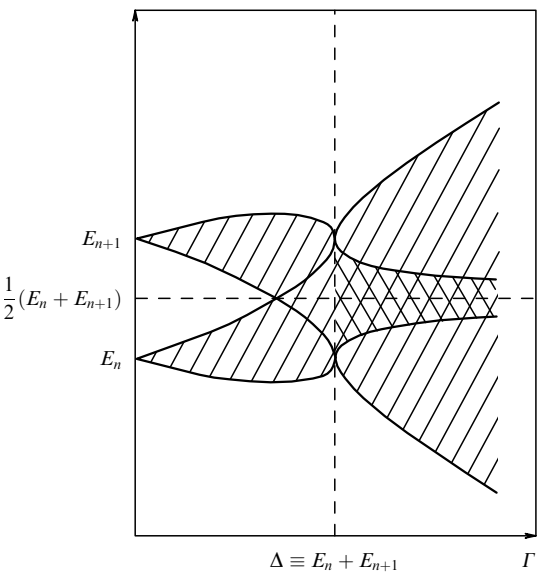

Figure 5. Broadened quasi-energy levels (bands) of the two-level system as functions of the ionization width Γ, calculated according to FGR and proportional to the intensity I.

To improve the model and make it more realistic, many Rydberg levels with different values of the principal quantum number n and different values of the orbital moment l were taken into account (Fig. 6). As a result, the dependences of the ionization probability on the fluence F and the pulse duration τ were obtained [12] (Fig. 7), which proved to be in good qualitative agreement with experimental curves (Fig. 2). Note that the experimental conditions [3] corresponded to a quite narrow region of parameters, shown in Fig. 7 by dashed lines. Therefore, it is obvious that experiment [3] do not yield information on the behavior of curves $w_{\mathrm{res}}(\varepsilon_0)$ in the region of strong fields $\varepsilon_0 \gg \omega^{5/3}$, although it suggests that stabilization exists for $\varepsilon_0 \geqslant \omega^{5/3}$.

The most recent variant of the model under study [16] also takes into account the possibility of gradually switching on and off the interaction of an atom with the laser pulse.

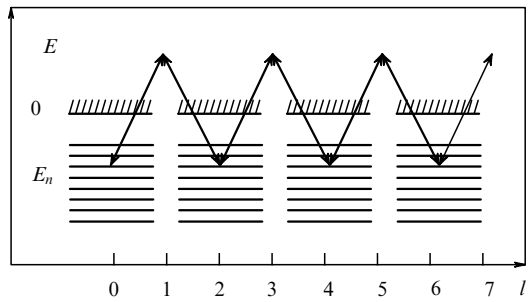

Figure 6. Transitions to the Rydberg states with relatively high values of orbital moment l.

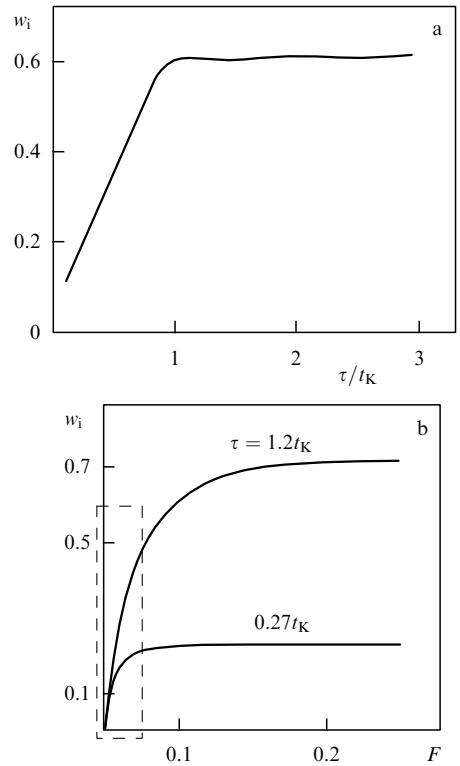

Figure 7. Theoretical dependences of the ionization probability of the Rydberg atom on (a) the pulse duration and (b) the fluence in a strong field [12].

Omitting details, we present here the most interesting results. The decay dynamics of an atom is characterized by the time-dependent total residual probability $w_{\mathrm{res}}(t)$ of finding it in discrete levels. In the case of a strong field, this dependence is shown in Fig. 8 for smooth and rectangular pulses $\varepsilon_0(t)$. The difference between these cases is quite noticeable, which is indicative of the drawback of the rectangular pulse model. Figure 8 shows that in the case of a smooth pulse envelope the ionization mainly occurs at the front and rear wings of the pulse, i.e., in a relatively weak field. At the central part of the pulse, where the field is strongest, the residual probability $w_{\mathrm{res}}(t)$ is almost constant. This is a direct manifestation of the stabilization of an atom in a strong field in the ionization dynamics. Figure 9 presents the dependences of the residual probability w_{res} at the instant $t = \tau$ of pulse termination on the parameter V proportional to $\varepsilon_0/\omega^{5/3}$ for different values of the pulse duration τ. Of most interest is the bottom curve corresponding to the longer pulse

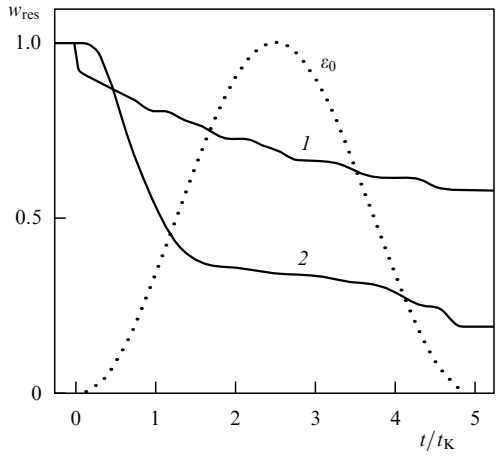

Figure 8. Theoretical dependences $w_{res}(t)$ of the 'non-ionization' probability on time for (*1*) a rectangular and (*2*) smooth pulses (the dotted curve is an envelope of the smooth pulse) [16].

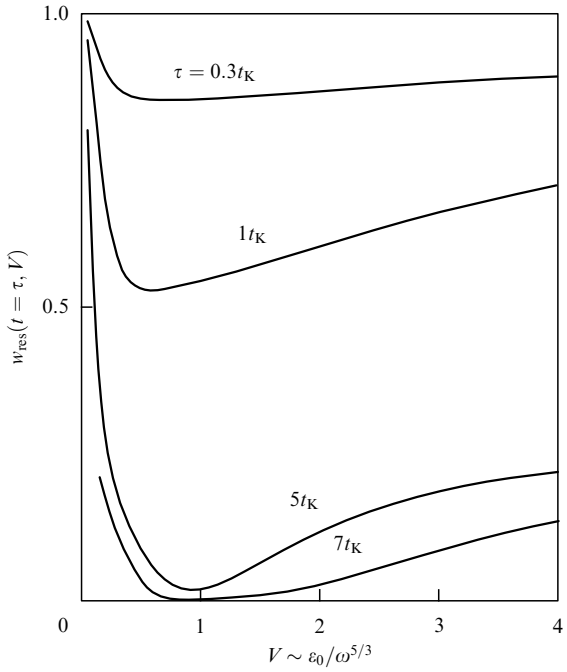

Figure 9. Residual probability of finding an atom in different bound states by the instant ($t = \tau$) of pulse termination as a function of the peak electric strength of the light pulse for different pulse durations τ [16].

duration ($\tau = 7t_K$). It shows that the atom can be stabilized when the pulse duration considerably exceeds the duration of the Kepler period $t_K = 2\pi n^3$. In this case, stabilization appears in fields that are somewhat stronger than in the case of short pulses ($\tau \leqslant t_K$).

Quasi-classical solution of the Schrödinger equation in a strong field. Although Eqns (5) for the probability amplitudes $C_n(t)$ and their solution appear quite reasonable, the above model itself is not rigorous. The main difficulties are encountered in the justification of the adiabatic elimination of the continuum. This procedure implies the use of a number of other approximations whose validity in a strong field is not obvious (the rotating wave approximation, the pole approximation, the flat continuum approximation, etc. [9]). In this connection, attempts have been made to develop a theory based on other models without using these approximations. One such approach is the attempt to use the quasi-classical approximation for the direct solution of the Schrödinger equation for the atomic electron in the presence of the time-dependent strong field [11, 14, 15]. Omitting the details, we present here the final expression for the total ionization probability of an atom by the pulse [14, 15]:

$$w_i = \int_0^1 dx \int_{-\infty}^{\infty} \frac{dt}{2t_K} \left[1 - J_0^2 \left(x \frac{2^{2/3} 3^{1/6} \Gamma(2/3) \varepsilon_0(t)}{\omega^{5/3}} \right) \right] . \quad (7)$$

This result was obtained without using the adiabatic elimination of the continuum and all the specific approximations that have been used in the model of equations for probability amplitudes $C_n(t)$. Expression (7) is valid for any shape of laser pulse envelope $\varepsilon_0(t)$. However, it was assumed in deriving (7) that the angular motion in a Rydberg atom is slow compared to the radial motion. This approximation has no rigorous justification. However, comparison with the exact numerical solution of the problem [17] directly confirms the validity of this approximation and Eqn (7). The results of calculations from Eqn (7) and obtained with the help of the exact numerical solution of the three-dimensional Schrödingr equation [17] for the model trapezoid-like function $\varepsilon_0(t)$ are presented in Fig. 10 (the solid and dotted lines, respectively). One can see that these results in fact completely coincide in the region of strong fields. Finally, it should be noted that the comparatively simple quasi-classical theory of photoionization of Rydberg atoms [14, 15] is restricted in the region of short pulses ($\tau < t_K$).

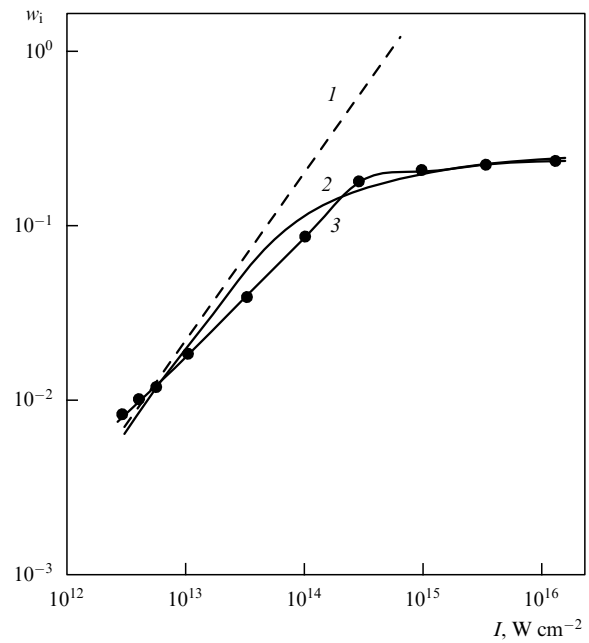

Figure 10. Ionization probability of the Rydberg atom during the pulse as a function of its peak intensity: (*1*) the perturbation theory, (*2*) the calculation from Eqn (7), and (*3*) the exact numerical solution [17].

5. Conclusions

In summary, we can state that while Kramers – Henneberger stabilization mainly remains a subject of theoretical discussions, the interference stabilization of Rydberg atoms has

been experimentally confirmed for the first time. However, many problems (for example, concerning the regimes of stabilization and ionization in a strong field) need experimental elucidation. As for the theory of interference stabilization, none of the models available and none of the approaches used are rigorous. Therefore, it is appropriate to continue extensive studies by different methods by comparing the results and finding the most reliable of them. It is also worthwhile extending the regions of parameters of exact numerical solutions and experimental studies. One of the most important new theoretical results is the possibility of stabilization of atoms in the case of long pulses ($\tau \gg t_K$) by increasing the field intensity. This result was obtained using a model of equations for the probability amplitudes $C_n(t)$ and has not yet been confirmed by other methods. In this connection, it is very important to extend the region of applicability of the solution based on the quasi-classical approach to pulses of long duration. It is also important to extend the range of parameters in which the problem can be exactly solved numerically.

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Hard X-ray radiation and fast particle generation using multiterawatt laser pulses

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1. Introduction

It is known that a laser plasma is an intense source of X-ray radiation and fast particles. This problem has been studied for nanosecond plasma in many papers [1]. At present, due to the advent of powerful picosecond and subpicosecond lasers, the possibility of generating X-rays and fast particles produced upon interaction of radiation from these lasers with solid targets in vacuum has appeared (see, for example, Ref. [2]). The physics of the interaction of ultrashort ($\tau \leqslant 1$ ps) powerful ($I > 10^{15}$ W cm^{-2}) laser pulses with solid targets considerably differs from the situation with long (nanosecond) pulses, because in this case the target material practically does not disperse, resulting in the formation of a high-temperature dense plasma in a small volume. This plasma is an intense source of X-ray radiation with the energy of quanta achieving several tens of MeV, and also of fast particles, first of all electrons, ions, positrons, and even muons [3].

In this article, we consider questions related to the generation of X-ray radiation and fast particles produced upon irradiation of solid targets by picosecond laser pulses with intensities exceeding 10^{16} W cm^{-2}. We will mainly analyze the results of experiments performed on our laser setups using such laser intensities.

2. Laser systems

To produce high intensities of focused laser radiation, we used two Nd:glass laser systems with the amplification and compression of a chirped pulse [4]. The laser system developed in the Institute of Laser Physics (Fig. 1) produced 1.5-ps pulses with an energy of up to 2.5 J and almost diffractive angular divergence of radiation. These pulses provided the intensity of radiation focused onto a spot of diameter about 10 μm up to 10^{18} W cm^{-2}. The high contrast of the pulse power, up to 10^7, was provided by a system of Pockels cells and electro-optical deflectors. Note that in a number of plasma experiments related, for example, to the optimization of the integrated yield of X-ray radiation, it is necessary to generate not only a single ultrashort pulse but also a train of pulses with a controllable time interval between them. Such a pulse train with the time interval controlled within 5 – 200 ps was obtained in this laser system by time multiplexing of the ultrashort pulse train from a master oscillator (MO) in a regenerative amplifier (RA) [5]. In this case, the pulse-repetition period was determined by the difference in bases of the resonators of the MO and RA, and the relative amplitude of a pulse in the train is controlled by electro-optical elements.

The laser system developed in the Institute of Complex Testing of Optical and Electronical Systems [6] differs from the above system by the method of obtaining a chirped pulse and the sizes of output amplifier cascades and diffraction gratings of the compressor. The high contrast of the laser pulse was provided by a multicascade electro-optical system based on Pockels cells. A block diagram of this system is presented in Fig. 2. It consists of four main parts: a driving laser system [6], a powerful amplifier cascade based on a 'Progress' amplifying channel [7], a compressor with holographic gratings, and a focusing system [8].

The driving laser produces chirped pulses of duration of the order of 300 ps and an energy of up to 1 J at a wavelength of 1053 nm. These pulses are then amplified in a powerful amplifier or compressed to 1.5 ps with the help of two 1700 lines mm^{-1}, 110×170 mm diffraction holographic gratings to perform plasma experiments with a power density of the order of 10^{17} W cm^{-2}.

A single pulse is separated and gated in a system of four Pockels cells, which are controlled by synchronized high-voltage electric pulse generators based on drift diodes with fast restoration of the reverse voltage [9]. These generators produce electric pulses of 1.5 – 5 ns duration at the half-height with amplitudes up to 15 kV.

The power of laser pulses was increased to 20 – 30 TW in a three-cascade amplifier channel containing three Nd:glass rod amplifiers with apertures 4.5, 6, and 8.5 cm and 30 cm in length [10].