Suppression of resonant multiphoton ionization via Rydberg states

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We discuss atomic stabilization in the process of resonance ionization via Rydberg states. We show that a strong resonance interaction between Rydberg states and a low-lying state substantially suppresses ionization. As a result, the population is trapped in the Rydberg states. Raman-type coupling of Rydberg states via the continuum is shown to be essential for the effect even when the bound-free interaction is weak.

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I. INTRODUCTION

One of the most interesting recent results in the physics of laser-atom interaction is the discovery of atomic stability against ionization in intense laser fields [1]. In this paper, we discuss one of the possible stabilization mechanisms. We show that resonance ionization via Rydberg states can be suppressed owing to strong resonant coupling of Rydberg states to the lower-lying bound state.

Several mechanisms of stabilization have recently been discussed [1–5]. Some are related to the stabilization of Rydberg states against one-photon ionization under different conditions [3–5]. The problem considered here is close to that studied in Ref. [3], and the relation between our results and those of Ref. [3] are discussed below. In this paper, we show that the mechanism we discuss, like other mechanisms of stabilization of Rydberg states, requires Raman-type coupling of the states via the common continuum, which leads to the destructive interference of bound-free transitions. We emphasize the role of strong resonant coupling and show that it can be responsible for this destructive interference.

In the field of a single-frequency laser pulse, Raman-type transitions between Rydberg levels are nonresonant and ionization suppression arises when the resonance detuning is compensated by large bound-free coupling, or, qualitatively, by an effective broadening of the Rydberg states due to strong bound-free transitions [4]. Alternatively, the detuning can also be compensated by the Fourier-limited bandwidth of an ultrashort pulse [5], in which case ionization may be suppressed when the bandwidth is much larger than the distance between neighboring Rydberg levels.

The suppression of ionization we discuss here is due to the strong interaction of the initial (e.g., ground) state with Rydberg states. It occurs even when the bound-free transitions are weak and the pulse duration is long. Nevertheless, coupling of Rydberg states via the continuum is central to this stabilization mechanism; stabilization disappears completely when this coupling is not taken into account. Qualitatively, an interaction with the ground state leads to the so-called effective *field broadening* [6] of Rydberg states which compensates for the detuning of Raman-type coupling via the continuum.

A pictorial scheme for the process is shown in Fig. 1. The initial state is either the ground or a lower excited s state of an atom. The laser field $\mathcal{E}(t) = \mathcal{E}_0 f(t)(\exp(-i\omega_L t)+\text{c.c.})$ is linearly polarized along the z axis. Absorption of one laser photon $\hbar\omega_L$ excites the atom to high Rydberg states of p series ($n \gg 1$, $l=1$); absorption of another photon ionizes the atom to the s or d continuum. We do not include the coupling of $l=2$ continuum states to $l=3$ Rydberg series (which is decoupled from the ground state). In other words, we neglect the migration of population from np Rydberg states towards $l=3$ and, possibly, $l=5, \ldots$ states. This approximation will be discussed below.

We call coupling of the initial state $|g\rangle$ and the Rydberg states $|n\rangle$ strong when

$$W_{gn} \equiv z_{gn}\mathcal{E}_0 > n^{-3}, \tag{1}$$

where z_{gn} is the transition matrix element between the states $|g\rangle$ and $|n\rangle$, and n^{-3} is the distance between adjacent Rydberg levels. Under the condition of Eq. (1) many

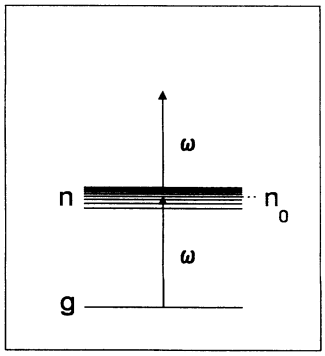

FIG. 1. Qualitative scheme of the process; n_0 is the Rydberg state closest to resonance.

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adjacent Rydberg levels resonantly interact with the initial state $|g\rangle$. On the other hand, we assume a weak bound-free coupling, that is, $\pi|V_{nE}|^2 \equiv \pi|z_{nE}\mathcal{E}_0|^2 \ll n^{-3}$, where $|E\rangle$ is the continuum state and $E=E_n+\omega$. Note that $|V_{nE}|^2$ includes the coupling to both s and d continuum states, $|V_{nE}|^2=|V_{nE}^{(s)}|^2+|V_{nE}^{(d)}|^2$.

It is well known [7] that for high Rydberg states the matrix elements z_{gn} and z_{nE} depend on the principal quantum number as $n^{-3/2}$: $z_{gn}=z_1 n^{-3/2}$, $z_{nE}=z_2 n^{-3/2}$, where z_1, z_2 are n independent. Therefore, Eq. (1) requires $\mathcal{E}_0 > 1/(n^{3/2}z_1)$. This condition becomes less severe with increasing n. On the other hand, the condition for weak bound-free coupling is n independent: $\mathcal{E}_0 \ll 1/(\pi z_2)$.

It is useful to give a quantitative estimate for the intensities required to satisfy the conditions of strong bound-bound and weak bound-free couplings. According to Ref. [7], an asymptotic form for a $2s$-np dipole matrix element in hydrogen is $z_{2s,np} \approx 4n^{-3/2}$. Hence, for $n \approx 50$ the condition of Eq. (1) is satisfied for intensities $I \geq 10^{10}$ W/cm^2. On the other hand, bound-free coupling remains weak until $I \sim 10^{13}$ W/cm^2. Note that this latter estimate is n independent for a given laser frequency, which is fixed by the initial state (e.g., $2s$).

Under the condition of Eq. (1) many Rydberg states are in resonance with the initial state and hence a Rydberg wave packet is excited when the field is turned on [9]. Upon passing the nucleus while moving along the Kepler orbit, the wave packet can proceed along three channels: (i) emit a photon, (ii) continue its motion along the Kepler orbit, and (iii) absorb a photon and ionize. As we show below, strong coupling to the initial state suppresses the third channel and favors the first two.

It is interesting to compare this quantitative picture, which is discussed in detail below, to the results of numerical simulations [3]. In Ref. [3] resonant two-photon ionization of the $2s$ state of hydrogen was studied, with one-photon resonance at not very high Rydberg states (typically $n \sim 10$), at laser intensities $I \geq 10^{13}$ W/cm^2. The spatial structure of the excited Rydberg wave packet suppressed further ionization; most of the wave function was localized far from the nucleus, where photon absorption is weak, and stabilization was attributed to this spatial structure [3(a)]. It is also important that partial return of population to the ground state was observed.

In our opinion, the spatial structure of the wave packet [3] was due to several factors: (i) short duration of the excitation pulse, (ii) Raman mixing of adjacent Rydberg states due to both *strong* coupling to the continuum and resonant coupling to the initial state, and (iii) fast spreading due to low values of principal quantum number n. In this work we analytically study the second factor, which is vitally important for stabilization. We also attempt to study it in a "clean" form. Therefore, we consider the case of large n of excited Rydberg states, so that the wave-packet spreading is slow. The pulse duration τ_L is long compared to the Kepler period $T_n=2\pi n^3$, so that the pulse bandwidth is narrow compared to the distance between neighboring Rydberg states. Finally, the bound-free coupling is weak, $|V_{nE}|n^3 \ll 1$, so that we can neglect above-threshold ionization and the migration of population to the states with high angular momenta (see below). Under these circumstances the only physical mechanism left to cause destructive interference of bound-free transition is strong interaction between Rydberg states and same initial state, which enhance the Raman coupling of Rydberg states via the continuum.

II. BASIC EQUATIONS AND APPROXIMATIONS

Let us focus on the mathematical formulation of the problem. In order to treat it analytically we need to make several approximations, which are discussed in succession. We also first consider the simplest case of an abrupt turn-on and turn-off of the laser pulse: $\mathcal{E}(t)=2\mathcal{E}_0 f(t)\cos\omega_L t$ with $f(t)=1$ at $0<t<\tau_L$ and $f(t)=0$ otherwise. Subsequently, the case of a smooth-pulse envelope is discussed.

Using the model of sharp turn-on and turn-off of the pulse requires some attention. It is well known that the Fourier spectrum of a step function is very wide. Therefore, one has to be careful about effects that can be introduced by the artificially broad spectrum. The model pulse we consider here is a difference between two step functions, and its Fourier spectrum is $\mathcal{E}(\Omega)/\mathcal{E}(\omega_L) \propto \sin(x)/x$, where $x \equiv (\Omega-\omega_L)\tau_L/2$. Obviously, the relevant value for one-photon transitions shown in Fig. 1 is laser intensity I. In our model the intensity spectrum is $I(\Omega)/I(\omega_L) = \sin^2[(\Omega-\omega_L)\tau_L/2]/[(\Omega-\omega_L)\tau_L/2]^2$, with the width $\Delta\Omega \sim 1/\tau_L$ determined only by the pulse duration τ_L. Therefore, for pulses with duration $\tau_L \gg 2\pi n^3$ we can neglect possible direct one-photon ionization from the initial ($2s$) state: the relative intensity at the required frequencies is about $(\tau_L/n^2)^2 \gg n^2$ times less than at the resonant frequency ω_L. Similarly, one can also ignore possible population of near-threshold continuum states due to one-photon transition from the Rydberg states: the laser frequency is much larger than their binding energy n^{-2}, and the spectral width is characterized by $\tau_L^{-1} \ll n^{-3}$. Hence, in the basic equations for essential states involved in the dynamics it is sufficient to take into account coupling of the initial state to the Rydberg states close to the resonance and resonant coupling of the Rydberg states to high states in the continuum. As for the nonresonant interaction of any given Rydberg state with its neighbors and near-threshold continuum states, it only contributes to the ac- Stark shift which is equal to $E_0^2/4\omega_L^2$ and is incorporated into the energies of Rydberg states.

As mentioned above, our model neglects the migration of population from $l=1$ to $l=3$ Rydberg states via common $l=2$ continuum. Indeed, $l=3$ states are decoupled from the ground state. Together with the condition $\pi|V_{nE}|^2 \ll n^{-3}$, it means that the migration can occur only if $l=1$ and 3 Rydberg states are degenerate, i.e., it is possible only in a hydrogen atom. Indeed, in complex atoms p and d Rydberg states have different quantum defects and their difference is on the order of unity. In other words, the detuning of the two-photon transition $np \rightarrow Ed \rightarrow nf$ is about n^{-3} and the two-photon coupling is much less. Strong resonant coupling to the initial states also cannot help—f states are decoupled from it.

Still, in hydrogen atom-population transfer to f Rydberg states is possible, owing to the accidental degeneracy of the states with different l. This problem was studied in Ref. [8] in some detail. It turned out that when exact dipole matrix elements of bound-free transitions are used, the migration $np \rightarrow Ed \rightarrow nf$ is very difficult. The reason is, in particular, the Bethe propensity rule for bound-free transitions. A detailed numerical study of this problem in the situation similar to ours was made by Huens and Piraux (Ref. [3c]), with the following conclusion: migration of population to $l=3$ Rydberg states becomes important only at very high intensities. Quantitatively, intensities above 10^{15} W/cm^2 (at $\omega=3.21$ eV) were required to obtain significant migration of population from p to f Rydberg states. At these intensities it is a very important effect that increases the stability of the atom [3(c)]. However, we are well below these intensities in our paper.

In our opinion, the basic origin of ineffective population transfer to the states with higher angular momentum is the ratio of laser frequency to the Kepler frequency: $\omega \gg n^{-3}$. Consequently, in moderate fields classical electron motion can be well separated into slow motion along the Kepler orbit and fast wiggling in the laser field. As a result, in the limit $\omega \gg n^{-3}$ the angular momentum turns out to be a conserved quantity. Only in a very strong field does the separation of motion into slow and fast fail, and only in this region does migration of population become important.

According to the above discussion, under the conditions considered in this paper the essential states in the problem are the initial state $|g\rangle$, resonant Rydberg states $|n\rangle$, and continuum states with energies close to $E_g+2\omega_L$ and $l=0,2$. Let $|n_0\rangle$ be the Rydberg state closest to the resonance, and choose its energy E_{n_0} as a zero-energy level (Fig. 1). Using the rotating-wave approximation (RWA), the probability amplitudes c_g, c_n (to find the atom in initial state $|g\rangle$ and Rydberg states $|n\rangle$, respectively) are given by the usual set of equations [5(b),9–12]

$$i\dot{c}_n = W_{ng}c_g + \Delta_n c_n - i\pi \sum_k V_{nE}V_{Ek}c_k \ , \tag{2a}$$

$$i\dot{c}_g = \sum_n W_{gn}c_n + \Delta_g c_g \ , \tag{2b}$$

where $\Delta_n = E_n - E_{n_0}$ and $\Delta_g = E_g + \omega_L - E_{n_0}$, Rydberg states are coupled via both s and d continuum, $V_{nE}V_{Ek} = V_{nE}^{(s)}V_{Ek}^{(s)} + V_{nE}^{(d)}V_{Ek}^{(d)}$, and the initial conditions are $c_g(0)=1$, $c_n(0)=0$.

Equation (2a) is not exact. Instead of two equations connecting the Rydberg-state amplitudes c_n to the continuum-state amplitudes c_E, we have only one equation where each Rydberg state c_n is coupled to all other Rydberg states c_k by the Raman-type transitions via the continuum [sum in Eq. (2a)]. Derivation of Eq. (2a) and the procedure of excluding the continuum is discussed in detail in the review [6] and Refs. [5(b),9(a),10], and the general approach is described by Faisal [11]. The most recent and very detailed discussion of the approximation equation (2a) can be found in Ref. [12]. In summary, the approximation equation (2a) is valid if the two requirements are met.

First, the excess of energy $E_g+2\omega_L \approx E_{n_0}+\omega_L$ over the ionization threshold must be sufficiently large, that is, much larger than the width of the zone of continuum states populated in the ionization process. This allows one to neglect the presence of ionization threshold in the integral over the energies of continuum states.

Second, the dependence of V_{nE} on E has to be smooth enough. Ideally, the matrix element V_{nE} should be E independent (flat continuum model) [5(b)]. This is never true for a real quantum system. However, as far as the RWA is valid and the width of the zone of populated continuum states is narrow compared to the excess of energy $E_{n_0}+\omega_L$ over the ionization threshold, the requirement to the function $V_{nE}=F_n(E)$ is much less severe. Namely, it is required that the continuum should be flat *within the width* Γ of the zone of populated continuum states. Physically, this result is transparent—as far as the RWA is valid, the only relevant states in the problem are those close to resonance.

As we see below, in our case the largest energy scale is given by $\Gamma \sim |W_{ng}|^2 n^3$, which determines the widths of the populated zones, both in the Rydberg spectrum and in the continuum. Obviously, we should require $\Gamma \ll n^{-2}$ to avoid population of near-threshold states due to resonant interaction with the initial state. Therefore, in our case the requirement to the dependence of V_{nE} on E is very easy to meet: the continuum should be flat within the width $\Gamma \sim |W_{ng}|^2 n^3$ near the energy $E_{n_0}+\omega_L$. The presence of the ionization threshold can be ignored as far as $\Gamma \ll E_{n_0}+\omega_L$, which is automatically fulfilled because $\Gamma \ll n^{-2}$ and $\omega_L \gg n^{-2}$.

We solve Eqs. (2), using the method of Laplace transformation, and that W_{gn}/V_{En} is n independent for $n \gg 1$. In the Laplace space, the following expression for $C_g(z)$, the Laplace transform of $c_g(t)$, can be obtained:

$$C_g(z) = \left[z + i\Delta_g + \frac{|W_{gn}|^2}{\pi |V_{nE}|^2} \frac{\phi(z)}{1+\phi(z)} \right]^{-1} ,$$

$$\phi(z) = \pi \sum_n \frac{|V_{nE}|^2}{z+i\Delta_n} \ . \tag{3}$$

Within the main assumption described above, i.e., sufficiently large excess of energy $E_{n_0}+\omega_L$ over the ionization threshold, the inverse Laplace transformation of $C_g(z)$ can be reduced to the inverse Fourier transformation with the change of variable $z=-i\epsilon$ [6,9(a),10,12]:

$$c_g(t) = \frac{1}{2\pi} \int_{-\infty}^{\infty} C_g(-i\epsilon)\exp(-i\epsilon t)d\epsilon \ . \tag{4}$$

The integral can be calculated using the residue theorem and is given by a sum of residues of $C_g(-i\epsilon)$ in the complex poles $\epsilon^{(n)}$, with the coefficients $\exp(-i\epsilon^{(n)}t)$. The ground-state amplitude must decay with time, and hence the complex poles $\epsilon^{(n)}$ must have a negative imaginary part, $\epsilon^{(n)} = \omega^{(n)} - i\gamma^{(n)}$. Then, completing the integration contour clockwise in the lower-half plane, we obtain

$$c_g(t)=-i\sum_n \mathrm{Res}[C_g(-i\epsilon^{(n)})]\exp(-i\epsilon^{(n)}t)\ . \quad (5)$$

From Eq. (5) it is clear that the poles $\epsilon^{(n)}$ determine the complex quasienergies of the dressed atom [11,12]. The real part $\omega^{(n)}$ determines the position of the quasienergy $E^{(n)}$ with respect to the above chosen zero-energy level E_{n_0}: $\omega^{(n)}=E^{(n)}-E_{n_0}$. The imaginary part $\gamma^{(n)}$ gives the decay rates of complex quasienergy states.

Furthermore, the residues $\mathrm{Res}[C_g(-i\epsilon^{(n)}]$ give the initial populations $w_0^{(n)}$ of complex quasienergy states. Indeed, by expressing the wave function as a superposition of the quasienergy states of the dressed atom, projecting it onto the initial state $|g\rangle$ and using the initial condition $c_g(t=0)=1$, we obtain

$$w_0^{(n)}=-i\,\mathrm{Res}[C_g(-i\epsilon^{(n)})]\ . \quad (6)$$

These populations decay with rates $2\gamma^{(n)}$. The ionization probability measured at the end of the pulse is

$$W_i(\tau_L)=1-\sum_n w_0^{(n)}\exp(-2\gamma^{(n)}\tau_L)\ . \quad (7)$$

Therefore, the problem of studying the ionization dynamics is reduced to the problem of finding the poles $\epsilon^{(n)}$ of $C_g(-i\epsilon)$, Eq. (3). The poles can be found using the semiclassical expression for $\phi(-i\epsilon)$ in Eq. (3) [13]:

$$\phi(-i\epsilon)=i\pi^2|V_{nE}|^2n^3\cot[\pi\nu(\epsilon)]\ ,$$

$$\nu(\epsilon)\equiv n_0/\sqrt{1-2\epsilon n_0^2}\ . \quad (8)$$

This expression is valid when $\nu(\epsilon)\gg 1$.

Equations (3) and (6)–(8) contain all the information we need to know about the ionization process. In order to extract this information we will first find the poles of $C_g(-i\epsilon)$, which will give us both the positions of the quasienergy levels and their widths. Then we will calculate the residues and thus find the ionization probability as a function of pulse duration.

III. RESULTS AND DISCUSSION

A. Spectrum of quasienergies

We first discuss the position of quasienergy levels $E^{(n)}$, and then their decay rates. Using Eqs. (3) and (8) we find that the position of the quasienergies $E^{(n)}=E_{n_0}+\omega^{(n)}$ depends on whether they are close to or far from the resonance; that is, whether $E^{(n)}-(E_g+\omega_L)$ is much less or much larger than the width of the resonance Γ, where Γ is given by

$$\Gamma=2\pi|W_{ng}|^2n^3\ . \quad (9)$$

Note that this expression coincides with a simple Fermi golden rule for the transitions to the quasicontinuum of states with density $\rho_n=n^3$. This is not surprising, as the "field broadening" W_{ng} [6] exceeds the level separation.

Close to the resonance, i.e., when $|E^{(n)}-(E_g+\omega_L)|\ll\Gamma$, we find that the quasienergy levels are localized almost in the middle between the field-free levels: $E^{(n)}\approx(E_n+E_{n+1})/2$. This result is similar to the case of strong bound-free coupling [4,10]. The difference is that now the reconstruction of the Rydberg spectrum is due to strong resonance interaction with the lower-lying isolated discrete state rather than due to interaction with the continuum. Another difference is that far from the resonance $[|E^{(n)}-(E_g+\omega_L)|\gg\Gamma]$ the spectrum recovers its field-free structure, that is, $E^{(n)}\approx E_n$.

We now discuss the ionization rates of complex quasienergies $E^{(n)}$. We show that strong resonance interaction with the initial state ($W_{ng}\gg n_0^{-3}$) suppresses ionization decay of complex quasienergy states. Using Eqs. (3) and (8) we find that $\gamma^{(n)}$ is determined by the following relation:

$$\gamma^{(n)}=\pi|V_{nE}|^2\frac{[E^{(n)}-(E_g+\omega_L)]^2}{\Gamma^2/4+[E^{(n)}-(E_g+\omega_L)]^2}\ . \quad (10)$$

Equation (10) has a form of standard Fermi golden rule ($\gamma_{\mathrm{FGR}}=\pi|V_{nE}|^2$) multiplied by an additional factor. From Eq. (10) it is immediately clear that the ionization decay of those states $E^{(n)}$, which are close to the resonance [i.e., with $|E^{(n)}-(E_g+\omega_L)|\ll\Gamma$], is strongly suppressed. Indeed, recalling that $E_g+\omega_L\approx E_{n_0}$, we can estimate the resonance detuning in Eq. (10) as $E^{(n)}-(E_g+\omega_L)\approx(n-n_0)n_0^{-3}$. Using this estimate in Eq. (10), we find that states with energies close to resonance $[|E^{(n)}-(E_g+\omega_L)|\ll\Gamma]$ ionize approximately $N^2/(n-n_0)^2$ times slower than they would normally do, with

$$N\equiv n_0^3\Gamma/2=\pi|W_{ng}|^2n^3n_0^3\gg 1\ , \quad (11)$$

the number of states in the half-width of resonance $\Gamma/2$. The closer the state to the resonance, the slower its decay. Far from the resonance ($|n-n_0|\gg N$) ionization is determined by the Fermi golden rule: $\gamma^{(n)}=\pi|V_{nE}|^2$.

It is strong resonance interaction with the initial state $|g\rangle$ that is responsible for the stabilization of the states $E^{(n)}$. However, if the Raman-type coupling via the continuum is not taken into account in Eq. (2a), the amplitude $C_g(z)$ will be determined by an equation different from Eq. (3), and one can easily show that the widths of the complex quasienergies will be $\gamma^{(n)}=\pi|V_{nE}|^2$, without the small factor $(n-n_0)^2/N^2\ll 1$. In sum, the stabilization mechanism we are discussing here is due to both Raman-type coupling of Rydberg states via the continuum and strong resonance interaction with the initial state, which compensates for the detunings of the Raman-type transitions.

The result of Eq. (10) has an interesting relation to an earlier result by Cardimona, Raymer, and Stroud [14]. If the laser frequency is tuned to a resonance with a *reconstructed Rydberg state,* that is, $E_g+\omega_L=E^{(n_0)}$ for some particular n_0, the corresponding width of *this particular* level is exactly equal to zero. This is exactly the effect discussed in Ref. [14]. No wonder that for high Rydberg states specific laser frequency corresponds to tuning approximately in the middle between Rydberg states n_0 and n_0+1. However, Eq. (10) contains much more information than this simple result. It says that *any* Rydberg level, which is sufficiently close to the resonance, is strongly

stabilized. This effect is due to strong resonance coupling and is frequency independent, that is, it occurs for any detuning between $E_g+\omega_L$ and a Rydberg state n_0. It also completely disappears if Γ becomes less than n^{-3}, and then only the effect described in Ref. [14] survives.

B. Ionization dynamics

In order to describe the ionization dynamics, we need to find initial populations $w_0^{(n)}$ of complex quasienergies. The probability of finding the atom in the quasienergy state $E^{(n)}$ at time t is $w^{(n)}(t)=w_0^{(n)}\exp(-2\gamma^{(n)}t)$, and the ionization probability is given by Eq. (7). As was discussed earlier, $w_0^{(n)}$ are given by Eq. (6), and the poles of $C_g(-i\epsilon)$ have already been found. Calculating the corresponding residues, we obtain

$$w_0^{(n)}=\frac{1}{\pi}\frac{\Gamma/2}{\Gamma^2/4+[E^{(n)}-(E_g+\omega_L)]^2} . \tag{12}$$

Qualitatively, the result of Eq. (12) is quite natural. Indeed, at $t \ll n_0^3$ the discrete structure of the Rydberg states is not resolved, and the transition from the ground state is similar to the bound-free transition. The transition rate $\Gamma=2\pi|W_{ng}|^2n^3$ is large ($\Gamma n_0^3 \gg 1$), and the ground state is totally depleted in a time much shorter than n_0^3. As a result, a wave packet with Lorentzian shape is created, similar to ionization. However, in our case the wave packet is bound and ionizes slowly, with the ionization rate of each state $E^{(n)}$ given by the corresponding $\gamma^{(n)}$. As $\gamma^{(n)}$ increases with increasing $|E^{(n)}-(E_g+\omega_L)|$ [see Eq. (10)], the decay of the wave packet begins at the wave packet's periphery and moves slowly to the center. In the absence of stabilization the decay rate is $2\gamma_{\rm FGR}=2\pi|V_{nE}|^2$, so it is convenient to use $\Theta=2\gamma_{\rm FGR}t$ as a time variable. Without stabilization the ionization decay of the wave packet would occur at $\Theta\sim 1$. In our case the ionization decay is much slower. Figure 2 shows the probability of finding the un-ionized atom at the moment t: $W(t)=1-W_i(t)$. The decay of $W(t)$ is nonexponential, and using Eqs. (10) and (12) one can show that $W(t)\propto 1/\sqrt{\Theta}$ for $\Theta>>1$. The dashed curve shows the ionization decay in the absence of stabilization, which is much faster.

So far the model has assumed an abrupt turn-on of the laser pulse. An optical pulse with a rise time τ_r shorter than the Kepler period $T_{n_0}=2\pi n_0^3$ for resonant Rydberg states and a long flat top is realistic [14]. Then, similar to the abrupt turn-on, a Rydberg wave packet will be created and, while moving along the Kepler orbit, it returns to the nucleus at $t\sim T_{n_0}$, the pulse intensity having already achieved its peak value. The only difference will be in the shape of the Rydberg wave packet, that is, in the initial populations of the quasienergy states. If $2\pi|W_{ng}|^2n^3\tau_r=\Gamma\tau_r \ll 1$, no significant excitation will occur at the pulse front and the Rydberg wave packet will be the same as in the case of abrupt turn-on. If $\Gamma\tau_r \gg 1$, the shape of the wave packet and initial populations $w_0^{(n)}$ will be determined by Γ, τ_r, and the pulse shape. The pulse should also be turned off much faster than in one Kepler period. Such a turn-off will not affect the wave packet moving far from the nucleus, where it spends most part of the Kepler period.

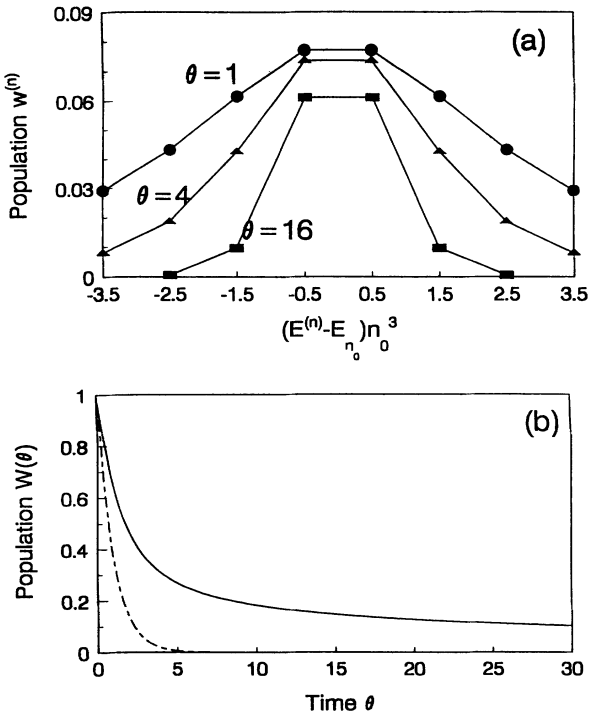

FIG. 2. Ionization decay of a model atom. $W(\Theta)$ is the residual population of all bound states, $W(\Theta)=1-W_i(\Theta)$. The number of states in a half-width of resonance [Eq. (11)] is $N=6$. Dashed curve—Fermi golden rule ionization.

Another important problem is the ponderomotive shift of the Rydberg levels. In the model it is time independent and is included in the energies E_n and resonance conditions. In the case of a smooth pulse it is time dependent and can shift the Rydberg states from the resonance. However, as long as $\omega_L \ll 1$, for *one-photon* resonance the shift, $\Delta E_{ac}=\mathcal{E}_0^2/4\omega_L^2$, is much smaller than the resonance width Γ. Indeed, semiclassical estimates for the matrix elements of transitions to and from Rydberg levels [4] give the ratio $\Delta E_{ac}/\Gamma\sim\omega_L^{4/3}\ll 1$. Therefore, resonance will not be affected by the Stark shift of Rydberg states.

IV. CONCLUSION

In summary, in this paper we have presented an analytical study of the process of multiphoton ionization under the conditions when the initial low-lying state is strongly coupled to a series of high Rydberg states. We have shown that for a long laser pulse $\tau_L > 2\pi n^3$ strong resonance interaction with high Rydberg states leads to the excitation of a Rydberg wave packet which is stable against further ionization. The stability of the wave packet is due to the coupling of many Rydberg states to the same initial state below, which gives rise to an efficient destructive interference of transitions to the continuum. Under the conditions of strong resonant coupling this effect is not sensitive to the laser frequency: the

laser can be tuned to any point between adjacent Rydberg levels. All Rydberg states sufficiently close to the resonance are strongly stabilized, and outside the resonance zone $\Gamma = 2\pi |W_{ng}|^2 n^3 \gg n^{-3}$ the original structure of the Rydberg states is restored.

Experimentally, the effect predicted in this paper can be observed by measuring the population remaining in the Rydberg states after the end of the pulse, as a function of pulse duration or pulse intensity. According to the calculations presented here, with increasing intensity the number of resonantly excited states N increases as $|W_{ng}|^2$. Hence, the population excited to each Rydberg state decreases as $1/N \propto 1/|W_{ng}|^2$. However, the ionization lifetime of each Rydberg state close to resonance will increase as $N^2 \propto |W_{ng}|^4$. Experiments could be done using ground-state Cs and a tunable dye-laser operating at a wavelength about 318 nm. The intensity, rise time, and pulse duration required depend on the principal quantum number of the Rydberg states to be excited. The characteristic time scale is determined by the Kepler period of excited Rydberg states; for states with $n \sim 70$ it is about 50 psec.

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[1] M. Pont, N. R. Walet, M. Gavrila, and C. W. McCurdy, Phys. Rev. Lett. **61**, 939 (1988); M. V. Fedorov and A. M. Movsesian, J. Phys. B **21**, L155 (1988); M. Yu. Ivanov, Bull. Acad. Sci. USSR, Phys. Ser. **52**, 117 (1988); Q. Su, J. H. Eberly, and J. Javanainen, Phys. Rev. Lett. **64**, 862 (1990); M. Dorr, R. Potvliege, and R. Shakeshaft, *ibid.* **64**, 2003 (1990); K. Kulander, K. Schafer, and J. Krause, *ibid.* **66**, 2601 (1991); R. R. Jones and P. H. Bucksbaum, *ibid.* **67**, 3215 (1991); for earliest references, see I. Gersten and M. H. Mittleman, J. Phys. B **9**, 2561 (1976); M. Gavrila and J. Z. Kaminski, Phys. Rev. Lett. **52**, 613 (1984).

[2] J. Groshmalicki, M. Lewenstein, and K. Rzazewski, Phys. Rev. Lett. **66**, 1038 (1991); M. Gajda, J. Groshmalicki, M. Lewenstein, and K. Rzazewski, Phys. Rev. A **46**, 1638 (1992); V. C. Reed, P. L. Knight, and K. Burnett, Phys. Rev. Lett. **67**, 1415 (1991); T. P. Grozdanov, P. S. Krstic, and M. H. Mittleman, Phys. Lett. A **149**, 144 (1990); H. R. Reiss, Phys. Rev. A **46**, 391 (1992).

[3] (a) K. Burnett, P. L. Knight, B. Piraux, and V. C. Reed, Phys. Rev. Lett. **66**, 301 (1991); (b) B. Piraux, E. Huens, and P. Knight, Phys. Rev. A **44**, 721 (1991); (c) E. Huens and B. Piraux, *ibid.* **47**, 1568 (1993).

[4] M. V. Fedorov and M. Yu.Ivanov, J. Opt. Soc. Am. B 7, 569 (1990).

[5] (a) J. Parker and C. R. Stroud, Phys. Rev. A **40**, 5651 (1989); (b) **41**, 1602 (1990); and (c) L. D. Noordam, H. Stapelfeldt, D. I. Duncan, and T. F. Gallagher, Phys. Rev. Lett. **68**, 1496 (1992).

[6] M. V. Fedorov and A. E. Kazakov, Progr. Quantum Electron **13**, 1 (1989).

[7] H. Bethe and E. Salpeter, *Quantum Mechanics of One- and Two-Electron Atoms,* 2nd ed (Rosetta, New York, 1977).

[8] R. Grobe, G. Leuchs, and K. Rzazewski, Phys. Rev. A **34**, 1188 (1986).

[9] (a) A. I. Andriushin, A. E. Kazakov, and M. V. Fedorov, Zh. Eksp. Teor. Fiz., **76**, 1907 (1979) [Sov. Phys. JETP **49**, 966 (1979)]; (b) E. Kyrola, J. Opt. Soc. Am. B **1**, 737 (1984); (c) G. Alber and P. Zoller, Phys. Rev. A **37**, 377 (1988); (d) V. M. Akulin and N. V. Karlov, *Intense Resonance Interactions in Quantum Electronics* (World Scientific, Singapore, 1992).

[10] M. V. Fedorov and A. M. Movsesian, J. Opt. Soc. Am. B **6**, 928 (1989); **6**, 1504 (1989).

[11] F. Faisal, *Theory of Multiphoton Processes* (Plenum, New York, 1987).

[12] M. V. Fedorov, Laser Phys. **3**, 219 (1993).

[13] I. J. Bersons, Zh. Eksp. Teor. Fiz. **83**, 1276 (1982) [Sov. Phys. JETP **56**, 731 (1982)].

[14] D. A. Cardimona, M. G. Raymer, and C. R. Stroud, Jr., J. Phys. B **15**, 55 (1982).

[15] J. P. Heritage, A. M. Weiner, and R. N. Thurston, Opt. Lett. **10**, 609 (1985); A. M. Weiner, J. P. Heritage, and R. N. Thurston, *ibid.* **11**, 153 (1986); R. N. Thurston, J. P. Heritage, A. M. Weiner, and W. J. Tomlinson, IEEE J. Quantum Electron. **QE-22**, 682 (1986).