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Rabi oscillations in a two-level atom moving in an open resonator with spherical mirrors

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***Abstract*.** **The dynamics of interaction of an electromagnetic field with a two-level atom uniformly moving in an open resonator is analysed. The probability of the radiative transition of an excited atom to a lower state is studied as a function of the slope of its trajectory relative to the resonator axis and of the detuning of the atomic transition frequency from the resonator mode frequency. The study was performed in the case of strong coupling by the method of dressed states. The complicated dependences of the transition probability on the detuning and the slope of the atomic trajectory are analysed for small angles ($\sim 1^{\circ}$) for typical experimental conditions. It is shown that under certain conditions, a field can be efficiently excited in the resonator by atoms whose transition frequency considerably differs from the field-mode frequency when the detuning is not close to the Doppler shift.**

Keywords: *Rabi frequency, interaction of a two-level atom with a field, transition probability*.

1. Introduction

The simplest model of the interaction of matter with an electromagnetic field is a two-level atom interacting with monochromatic radiation. Despite its obvious simplicity, this model can explain and predict a number of effects studied by quantum electronics, both qualitatively and quantitatively. In this model, the interaction of an electromagnetic field with an atom in a resonator is characterised by the coupling constant

$$g_0 = d\left(\frac{2\pi\omega_{\mathrm{a}}}{\hbar V}\right)^{1/2},$$

where d is the projection of a transition dipole moment of an atom on the direction of polarisation of the field mode under study; ω_{a} is the transition frequency for a two-level atom; and V is the resonator volume.

The model of interaction with a monochromatic field (single-mode approximation) can be applied when the interaction of the atom with the continuous spectrum of the electromagnetic field of free space surrounding the resonator can be neglected under some physical conditions, because this field causes spontaneous radiation of the atom. If we assume that the rate of spontaneous emission of the atom is γ and the emission loss rate in the resonator is k, the atom will emit in two regimes, when $g_0 \gg \gamma,\ \kappa$ (strong coupling) or $g_0 \ll \gamma,\ \kappa$ (weak coupling). The strong coupling regime is characterised by the oscillatory dynamics of the field in the resonator and of the atomic state, whereas the weak coupling regime features the exponential decay. In this paper, the strong coupling of an atom with the single-mode resonator field is considered.

It is known that the transition frequency of a moving atom increases or decreases, depending on the mutual orientation of the velocity vector of the atom and the wave vector of the field. In paper [1], the movement of an atom was analysed in the field of a standing wave of a resonator with flat mirrors, as well as a particular case of its movement transverse to the fundamental mode of a resonator with spherical mirrors. The dependence of the type of oscillations of the atom on its velocity (in the case of the resonator with plane-parallel mirrors, on the Doppler shift), the coupling constant, and the detuning of the transition frequency from the mode frequency was studied.

A series of recent experiments with micromasers [2 – 9] have demonstrated the possibility of obtaining the strong coupling regime by using a cylindrical resonator and the Rydberg atoms. The strong coupling regime was observed in a number of experiments when Rydberg atoms propagated through an open resonator [10 – 14] in which the TEM_{n00} mode was excited. The theory that was used so far for the interpretation of such experiments neglects the effects related to the movement of atoms, which, in our opinion, can result in a substantial discrepancy with experiments under certain conditions.

In this paper, the dynamics of the interaction of an atom with the field in an open resonator is analysed in the case when the atom propagates through the resonator at an arbitrary angle to its axis. The calculations performed within the framework of the theory developed in this paper demonstrate an extremely high sensitivity of the dynamics of the field and atomic state to the relation between the parameters of the atomic movement (the velocity and direction of the movement) and the coupling parameter, as well as to the detuning of the atomic transition frequency from the resonator frequency.

It is shown that the motion of atoms should be taken into account in the analysis of current experimental studies [9 – 15].

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2. Transition probability

We represent the Hamiltonian of a nonconservative system consisting of a two-level uniformly moving atom and a mode of the quantisied electromagnetic field of a resonator with frequency ω_c in the form

$$H(t) = \hbar\omega_c a^+ a + \frac{1}{2}\hbar\omega_a \sigma^z + \hbar g(t) B, \quad (1)$$

where the operator

$$B = a^+ b + b^+ a \quad (2)$$

of interaction of the field with the atom is written in the rotating wave approximation in terms of the operators of transitions between the upper $|\uparrow\rangle$ and the lower $|\downarrow\rangle$ states of the atom: $b = |\downarrow\rangle\langle\uparrow|$, $b^+ = |\uparrow\rangle\langle\downarrow|$; $\sigma^z = b^+ b - b b^+$ is the inversion operator.

Consider the case when a moving atom intersects an open resonator with spherical mirrors through its centre at an angle of ϕ to the transverse axis x of the resonator. Assuming that the movement of the atom along this axis passing through the resonator centre is described by the expression $x(t) = \upsilon_x t$, we place the origin of coordinates ($x = 0$) at the distance $L/2$ from the resonator centre. In this case, the dependence of the coupling parameter (1) on time has the form

$$g(t) = g_0 \cos\left[k_c\left(\upsilon_z t - \frac{L}{2}\frac{\upsilon_x}{\upsilon_z}\right)\right] \exp\left[-\left(\frac{\upsilon_x t - L/2}{w_0}\right)^2\right],$$

$$g_0 = d\left(\frac{2\pi\omega_a}{\hbar V}\right)^{1/2}, \quad (3)$$

where $k_c = \omega_c / c$ is the wave vector of the resonator mode and w_0 is the radius of the waist of the Gaussian distribution of the fundamental resonator mode. One can see that the coupling parameter depends on the spatial distribution of the field in the resonator and the velocity vector $\boldsymbol{\upsilon} = \upsilon\{\sin\phi, 0, \cos\phi\}$.

In the case of the uniform movement of atoms considered here, the coupling parameter can be written in the form

$$g(t) = g_0 \cos(\Omega_D t + \Delta\Omega t - \varphi) \exp\left[-\left(\frac{\upsilon_x t - L/2}{w_0}\right)^2\right], \quad (4)$$

$$\Omega_D = \frac{\upsilon_z}{c}\Omega_a, \quad \Delta\Omega = \frac{\upsilon_z}{c}\Delta\omega, \quad \varphi = \frac{L}{2}\frac{\upsilon_x}{\upsilon_z}, \quad (5)$$

where Ω_D is the Doppler shift of the atomic transition frequency and $\Delta\omega = \omega_c - \omega_a$ is the detuning of the atomic transition frequency from the resonator-mode frequency. It is also assumed that the field mode is not degenerate because the distance between the resonator mirrors does not equal exactly to the radius of curvature of the mirrors, which have in turn weak ellipticity.

The time evolution of the atom interacting with the field is described in the Schrödinger representation by the propagator (the evolution operator)

$$U(t, t_0) = 1 - \frac{i}{\hbar}\int_{t_0}^{t} H(t')U(t', t_0)\,dt'; \quad (6)$$

and the state vector at the instant of time t has the form

$$|\psi(t)\rangle = U(t, t_0)|\psi(t_0)\rangle. \quad (7)$$

Let us divide a finite time interval $[t_0, t]$ into a great number $M \gg 1$ of small segments Δt and assume that the time dependence of the Hamiltonian of the system can be neglected. Then, the time evolution of the system in each of the segments is described by the propagator

$$|\psi(t_j)\rangle = \exp\left[-\frac{i}{\hbar}H(t_j)\Delta t\right]\psi(t_{j-1})\rangle,$$

$$t_j = t_{j-1} + \Delta t, \quad j = 1, \ldots, M, \quad (8)$$

and the state vector of the system at the instant of time t is

$$|\psi(t)\rangle = \exp\left[-\frac{i}{\hbar}H(t_M)\Delta t\right]\exp\left[-\frac{i}{\hbar}H(t_{M-1})\Delta t\right]$$

$$\ldots \exp\left[-\frac{i}{\hbar}H(t_1)\Delta t\right]|\psi(t_0)\rangle, \quad \Delta t = \frac{t}{M}. \quad (9)$$

Below, we will consider the case when the atom is in an excited state at the initial instant of time $t_0 = 0$, while the field is in a vacuum state, i.e., $|\psi(t_0 = 0)\rangle = |0, \uparrow\rangle$.

By using (1), (2), (4), (8), and (9), we obtain the recurrent expression [1]

$$|\psi(t = t_M)\rangle = A_{+,M}|+, 0\rangle_{t=t_M} + A_{-,M}|-, 0\rangle_{t=t_M}, \quad (10)$$

for the state vector of the system at the instant of time t, where $A_{\pm,M}$ are calculated using the recurrent relations

$$A_{+,t_j} = \exp\left[-i\Omega_{+,0}(t_j)\Delta t\right]\left[\cos(\Delta\theta_{0,j})A_{+,t_{j-1}}\right.$$

$$\left.+ \sin(\Delta\theta_{0,j})A_{-,t_{j-1}}\right],$$

$$A_{-,t_j} = \exp\left[-i\Omega_{-,0}(t_j)\Delta t\right]\left[-\sin(\Delta\theta_{0,j})A_{+,t_{j-1}}\right. \quad (11)$$

$$\left.+ \cos(\Delta\theta_{0,j})A_{-,t_{j-1}}\right],$$

$$\Delta\theta_{0,j} \equiv \theta_0(t_j) - \theta_0(t_{j-1})$$

for any $0 < j \leqslant M$ and

$$A_{+,t_1} = \sin\left[\theta_0(\Delta t)\right]\exp\left[-i\Omega_{+,0}(\Delta t)\Delta t\right],$$

$$A_{-,t_1} = \cos\left[\theta_0(\Delta t)\right]\exp\left[-i\Omega_{-,0}(\Delta t)\Delta t\right]. \quad (12)$$

The function $\theta_0(t)$ appearing in (11) and (12) has the form

$$\theta_0(t) = \frac{1}{2}\arctan\left[\frac{2g(t)}{\Delta\omega}\right]. \quad (13)$$

The dynamic Rabi frequency in vacuum is described by the expression [1]

$$\Omega_{\pm,0}(t) = \pm\left[\left(\frac{\Delta\omega}{2}\right)^2 + g^2(t)\right]^{1/2}. \quad (14)$$

The probability of the atomic transition to a lower state accompanied by emission of a photon into the resonator mode is

$$P_\downarrow(t) = |\langle 1\downarrow|\psi(t)\rangle|^2 = |\cos(\theta_0(t))A_{+,t} - \sin(\theta_0(t))A_{-,t}|^2. \quad (15)$$

We assume in our calculations that the de Broglie wavelength $\lambda = h/(mv)$ (where m is the atom mass) is far shorter than the wavelength of the resonator mode $\lambda_c = 2\pi/k_c$, i.e., the mass of the atom and (or) its velocity is high. The assumption about the uniformity of the movement of the atom, i.e., the neglect of the recoil momentum upon emission of the photon requires that the binding energy $\hbar g_0$ should be small compared to the kinetic energy $mv^2/2$ of the atom. The applicability of the classical approximation to the problem of the interaction of an atom moving in a resonator is analysed in detail, for example, in paper [16].

We also assume that the Q factor of the resonator is high enough for the time of interaction of the atom with the field to be much shorter than κ^{-1}, where κ is the rate of field dissipation due to losses at the resonator mirrors. The conditions considered here were realised, in particular, in papers [3, 4, 6, 14], where the authors used the usual Jaynes–Cummings model.

The iteration formula for the evolution operator (6) has, in the interaction representation, the form

$$U(t,t_0) = \sum_{n=0}^{\infty}\frac{(-\mathrm{i})^n}{n!}\int_{t_0}^{t}\mathrm{d}t_n\int_{t_0}^{t}\mathrm{d}t_{n-1}\ldots\int_{t_0}^{t}\mathrm{d}t_1$$

$$\times T[V(t_n)V(t_{n-1})\ldots V(t_1)], \quad (16)$$

$$U(t,t_0) = T\left\{\exp\left[-\mathrm{i}\int_{t_0}^{t}V(t')\mathrm{d}t'\right]\right\}, \quad (17)$$

where

$$V(t) = -\frac{\Delta\omega}{2}\sigma^z + g(t)B. \quad (18)$$

The symbol T in (16) and (17) denotes the time ordering of the operators in the integrands. Because the operator $V(t)$ commutes with itself for any $t_1 \neq t_2$, there is no need to perform the time ordering in (17) in our case.

As follows from (17), the wave function of the system can be found for any instant of time t using the relation

$$|\psi(t)\rangle = \left\{1 + \sum_{n=0}^{\infty}(-\mathrm{i})^n\frac{1}{n!}\left[\left[\int_0^t \mathrm{d}t'V(t')\right]^n\right\}|0,\uparrow\rangle. \quad (19)$$

We find from (19) that the transition probability can be written as a series

$$P_\downarrow(t) = \left|I(t)\sum_{n=0}^{\infty}(-1)^n\frac{[I^2(t)+\delta^2(t)]^n}{(2n+1)!}\right|^2, \quad (20)$$

where

$$I(t) \equiv \frac{g_0}{k_c v\sin\phi}\int_{\alpha}^{\beta(t)}\exp(-c\tau^2)\cos\tau\mathrm{d}\tau; \quad (21)$$

$$\alpha \equiv -\frac{k_c L}{2}\tan\phi;\ \tan\phi = \frac{v_z}{v_x};\ \beta(t) \equiv k_c\left(vt\sin\phi - \frac{L}{2}\tan\phi\right);$$

$$c \equiv \left(\frac{\cot\phi}{k_c w_0}\right)^2;\ \delta(t) \equiv -\frac{\Delta\omega}{2}t.$$

In the case of the exact resonance $\Delta\omega = 0$, we obtain from (20) and (21)

$$P_\downarrow(t) = \sin^2[I(t)]. \quad (22)$$

Expression (20), which allows one to calculate easily the transition probability in the case of a small detuning of the atomic transition frequency from the mode frequency, becomes inconvenient at large detunings when the alternating series in (20) converges very slowly. In this case, the transition probability can be easily calculated from expression (15).

Let us find now from (22) the probability of emission of a photon by an atom moving through a resonator, i.e., the value of $P_\downarrow(t)$ for $t = L/v$ under the condition that $L \gg 2w_0$. We will consider the case of small detunings $\Delta\omega \ll g_0$. One can easily see that

$$P_\downarrow(t=\infty) = \sin^2\left\{\sqrt{\pi}\frac{g_0 w_0}{v\cos\phi}\exp\left[-\left(\frac{w_0\omega_c}{2c_L}\tan\phi\right)^2\right]\right\}, \quad (23)$$

where c_L is the velocity of light.

It follows from (23) that for

$$\frac{g_0 w_0}{v\cos\phi}\sqrt{\pi} \ll 1,\ v \gg \frac{g_0 w_0}{\cos\phi}\sqrt{\pi}$$

the energy exchange between the excited atom and the resonator is impossible, and the states of the atom and field remain unchanged during the entire transit time of the atom in the resonator. The angle between the direction of the velocity vector of the atom and the transverse axis z of the resonator during the passage of the atom through the resonator centre is $0 \leqslant \phi < \phi_{max}$, where $\phi_{max} = \arctan\{\mu[\nu - 1 + (1-\mu^{-2})^{1/2}]\}$; $\mu \equiv R/A$; R is the radius of curvature of resonator mirrors; $2A$ is the diameter of the mirrors; $\nu \equiv D/(2R)$; and D is the distance between the mirrors. The radius of the waist of the Gaussian mode of the resonator with the geometry close to the confocal one can be found from the expression

$$w_0 = \left[\frac{2R}{k_c}\left(\frac{D}{2R-D}\right)^{1/2}\right]^{1/2}.$$

3. Rabi oscillations of a moving atom

The parameters of an almost confocal resonator used in the calculations performed within the framework of the theory presented above were similar to those used in paper [10], where Rabi oscillations of a Rydberg atom leaving the resonator were directly observed by varying the time of interaction (the atom velocity) between the atom and field. The radius of curvature of almost spherical mirrors was $R = 4$ cm, the distance between them was $D = 2.75$ cm, and the diameter of the mirrors was $2A = 4$ cm (see also [11, 12]). The degeneration of the TEM_{900} mode used in the experiment was removed by making the resonator mirrors weakly elliptic. The fundamental-mode frequency was $\omega_c = 3.21\times10^{11}$ s^{-1} and the radius of the waist of the Gaussian distribution of the field was $w_0 = 0.596$ cm.

Consider the atoms passing through the resonator centre at an arbitrary angle of their rectilinear trajectories to the x axis ($0 \leqslant \phi < \phi_{\max}$). Note that the trajectories of the atom inside the resonator passing through its centre do not exhaust all the possible trajectories of the rectilinear motion. However, taking into account that under our conditions, $\phi_{\max} \approx \pi/16 = 11.25°$, which is a small value, the trajectories that do not pass through the resonator centre have much smaller slopes relative to the resonator axis because a source of atoms is removed from the resonator by the distance that greatly exceeds the resonator size. For this reason, no calculations for such trajectories were performed.

Fig. 1 shows the time dependence of the transition probability $P_{\downarrow}(t)$ for $\phi = 6°$ for the initially excited atom and the vacuum state of the resonator. As follows from (23), for $\phi = 0$, the absorption resonance ($P_{\downarrow}(\infty) = 1$) is possible for $g_0 w_0/(v\sqrt{\pi}) = n + 1/2$ and the transmission resonance ($P_{\downarrow}(\infty) = 0$) for $g_0 w_0/(v\sqrt{\pi}) = n$, where n is an integer. In the first case, a photon is captured by the resonator, while the atom in the lower state leaves the resonator. In the second case, the state of the atom passing through the resonator does not change (the resonator is transparent). The parameters v, g_0, and $\Delta\omega$ given in Fig. 1 correspond to the capture (absorption) resonance for $\phi = 0$ (i.e., $P_{\downarrow}(\infty) = 1$). It follows from Fig. 1 that the probability of the photon capture sharply decreases even at the small slope $\phi = 6°$.

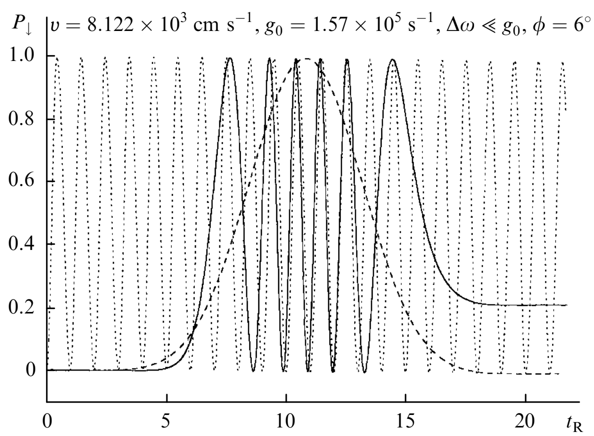

Figure 1. Probability $P(\downarrow)$ of the radiative transition in the excited atom to the lower state during its movement in an almost confocal resonator as a function of the dimensionless time $t_R = \Omega_R t/\pi$ (solid curve). The dashed curve shows the dependence $\cos\{k_c[v\cos\phi t_R - (l/2)\cot\phi]\} \times \exp\{-[(v\sin\phi t_R - (L/2))/w_0]^2\}$, and the dotted curve shows Rabi oscillations of an atom at rest.

In the case of a finite detuning $\Delta\omega \to g_0$, the resonator becomes transparent for $\phi = 0$ at any velocity. The shape of oscillations of the probability or of the average number of photons proves to be very sensitive to the magnitude of detuning $\Delta\omega = g_0$ even at its small, but comparable with the coupling constant g_0, values.

Fig. 2 shows the dynamics of the transition probability for the nonzero slope of the velocity vector of the atom relative to the resonator axis for the detuning $\Delta\omega = 0.2g_0$. One can see from Fig. 2 that the probability of the escape of the atom in the lower state from the resonator depends strongly and nonmonotonically on the angle ϕ. Fig. 3 shows the dependence of $P_{\downarrow}(\infty)$ on the angle ϕ, which was obtained from expression (23) for three velocities of the atom. One can see that this dependence becomes very strong with decreasing velocity. At the same time, when the trajectory of the atom is not parallel to the x axis, the transition probability strongly changes even at high velocities. Note that at some angles the probability of emission of a photon into the resonator proves to be independent of the atom velocity. Such separated directions of the movement of the atom inside the resonator under study correspond to angles $\phi \approx 3°$ and $\phi \approx 8.4°$.

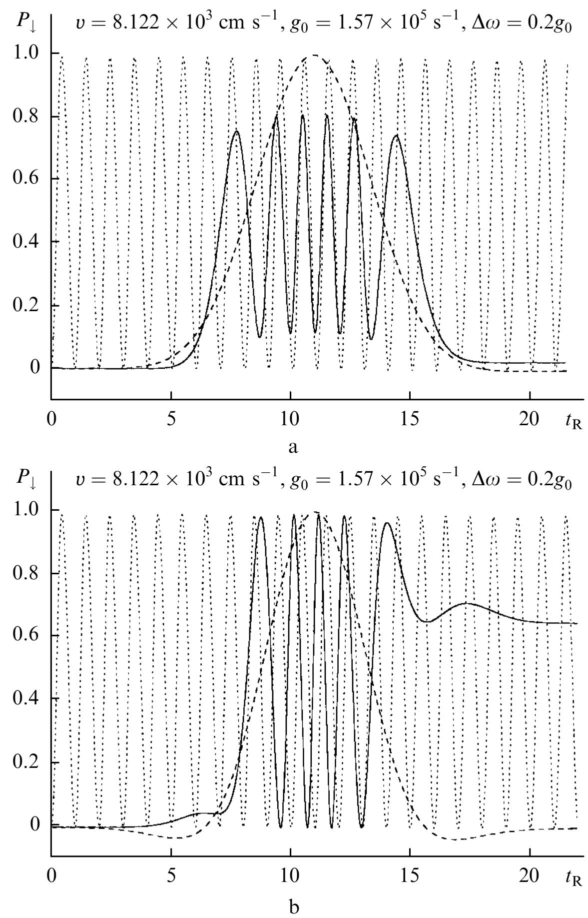

Figure 2. Same as in Fig. 1 for the angles between the atom trajectory and the resonator axis $\phi = 6°$ (a) and $\phi = 11.25°$ (b) but for $\Delta\omega = 0.2g_0$.

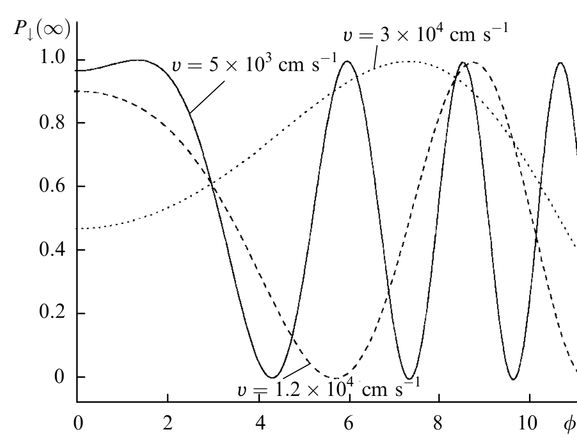

Figure 3. Dependences of $P_{\downarrow}(\infty)$ on the angle ϕ between the trajectory of the atom and the resonator axis for $\Delta\omega = 0.01g_0$ and different velocities v.

Fig. 4 shows the dependence of $P_\downarrow(\infty)$ on the detuning $\Delta\omega$ for different angles ϕ. One can see that for $\phi \neq 0$, the oscillations of the probability $P_\downarrow(\infty)$ exhibit the dependence on $\Delta\omega$ that is typical for the Rabi oscillations of the moving atom [1]. As the Doppler shift Ω_D increases with increasing angle ϕ, the compensation of the detuning of the transition frequency from the resonator-mode frequency occurs, and the transition probability increases at $\Delta\omega \sim \Omega_D$. Note that the maximum of the curve corresponding to $\phi = 11^\circ$ is observed when $2.3\Omega_D = \Delta\omega$, whereas for the resonance condition $\Delta\omega = \Omega_D$ [1] and a smaller detuning, the transition probability becomes substantially lower.

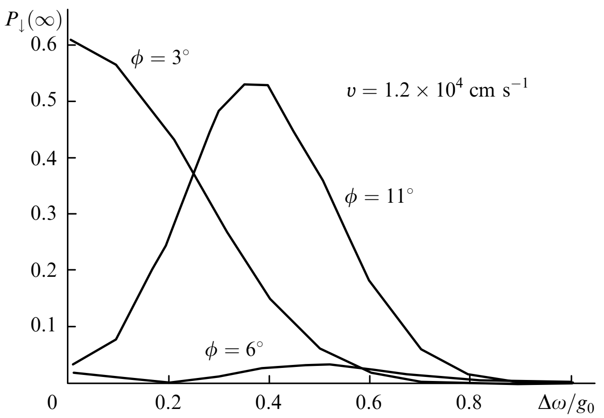

Figure 4. Dependences of $P_\downarrow(\infty)$ on the detuning $\Delta\omega$ for different slopes ϕ and $v = 1.2 \times 10^4$ cm s^{-1}.

The property found in this calculation leads to the following conclusion. When an excited atom with the transition frequency that substantially differs from the mode frequency moves at a certain angle to the resonator axis, the energy of the atom can efficiently transfer to the resonator mode at $|\Delta\omega| \neq \Omega_D$. Note that an atom with the transition frequency detuned from the mode frequency cannot excite the field in the resonator when it moves along the resonator axis [1], i.e., in the case of a common configuration used, for example, in the Ramsey spectroscopy [12] and in microlaser experiments [15].

4. Conclusions

The calculations performed in this paper showed that the movement of an atom plays an important role in the theoretical interpretation of experiments with atoms moving in a resonator (in particular, in micromaser and microlaser configurations), as well as in spectroscopic Ramsey measurements.

In the case of a small detuning of the atomic transition frequency from the mode frequency $|\Delta\omega| \ll g_0$, the average number of photons in the resonator after the propagation of the atom through the centre of the open resonator with spherical mirrors exhibits oscillations depending on the atom velocity. When $g_0 w_0/(v\sqrt{\pi}) = n + 1/2$, the atom transfers a resonance photon to the resonator with the unit probability. If $g_0 w_0/(v\sqrt{\pi}) = n$, the atom leaves the resonator in the initial excited state, whereas the resonator field remains unchanged. For $|\Delta\omega| \sim g_0$, the average number of photons in the resonator decreases with increasing detuning, and the probability of the field excitation in the resonator is close to zero already at $|\Delta\omega| \approx g_0$. When the atom moves in the resonator at an angle to its axis, the probability of excitation of the resonator by the nonresonance atom exhibits the nonmonotonic dependence on the slope even at small values of the slope.

As the slope ϕ and the atom velocity v increase, i.e., when the Doppler shift Ω_D increases, the efficiency of excitation of the field in the resonator by the atom with the detuning $\Delta\omega$ strongly increases. In this case, the detuning can be large ($|\Delta\omega| \sim \Omega_D$). Depending on the parameters of the system, the excitation probability proves to be maximal for $|\Delta\omega| > \Omega_D$, i.e., in the absence of a complete compensation for the detuning by the Doppler shift. Therefore, the efficient compensation for the detuning is achieved in the given configuration under special conditions (for a certain angle between the atom trajectory and the resonator axis). The excitation probability is high at a certain detuning for each value of the angle ϕ.

For small detunings ($|\Delta\omega| \ll g_0$), the probability of deexcitation of the atom and excitation of the resonator mode after the propagation of the atom is extremely sensitive to the value of the angle ϕ, especially, at low velocities of the atom. It has been also found that the transition probability is independent of the atom velocity for some angles ϕ. These angles are determined by the resonator geometry and the coupling parameter of the atom with the field.

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