Momentum transfer for an optical transition in a prepared two-level atom

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We consider the interaction of a traveling optical wave with a resonant two-level atom, which is initially in its general superposition state, i.e., in the superposition state of the ground and excited internal energy levels and simultaneously has mutually different momentum distributions, corresponding to each of these internal energy levels. We show that during the interaction, the atom periodically gets large-scale changes in these momentum distributions. As a consequence, a portion of the atomic momentum of the states, corresponding to each internal energy level, gets large-scale changes too, much more than the momentum $\hbar k$ of an absorbed/emitted photon. The amount of momentum that a photon transfers between the atomic internal energy levels is in general more than its own momentum $\hbar k$. The special case is discussed in which the atom's preliminary superposition state is created as a result of interaction of the atom with the resonant standing wave. Also it is pointed out that in appropriate time intervals the mentioned phenomenon can be presented as a transformation of the resonant Kapitza-Dirac splitting of atomic states into the Stern-Gerlach-type splitting.

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I. INTRODUCTION

When an atom interacts with a resonant traveling wave, the atomic total momentum changes in units of one photon momentum $\hbar k$. What can be said about distributions and mean values of momentum for translational states connected with each atomic level? The answer is well-known and is trivial, if the atom before the interaction is on one of the internal energy levels: the momentum distribution on the other level gets shifted by $\hbar k$ and the mean value of momentum gets shifted also by $\hbar k$; $\overline{p_e}=\overline{p_g}+\hbar k$, where $\overline{p_g}$ and $\overline{p_e}$ are mean values of momentum, corresponding to ground and excited internal levels consequently [one-dimensional (1D) case]. Therefore, a photon, during an absorption or emissions, transfers between the atomic internal energy levels an amount of momentum Δp just equal to its own momentum $\hbar k$. What would we have in a general case, that is, when the atom before the interaction with the traveling wave is in the superposition state of ground and excited levels with mutually different momentum distributions there, that is, when the ''two-level atom'' is in its most general superposition state? In the following just this question will be elucidated, rather in general form and for the important special case. It will be shown that in general already $\Delta p \neq \hbar k$, that is, a photon being absorbed or emitted by an atom, transfers between internal atomic energy levels the amount of momentum not necessarily equal to the photon's own momentum; moreover, this amount may even greatly exceed the photon's own momentum.

In Sec. II it will be shown that due to the interaction with the traveling wave the essential momentum distributions are possible in corresponding atomic states to each internal energy level. These redistributions are a direct consequence of interference between the amplitudes of translational states of the ground internal energy level and the excited one. Further we present conditions when the redistributions are large-scale and respective variations of mean momentum per each internal energy level greatly exceed one photon momentum $\hbar k$.

In Sec. III we will discuss from the practical point of view a very important case, when the preliminary superposition state of the atom is realized by the coherent diffraction of the atom in the field of a resonant standing wave, which is often being referred to as the resonant Kapitza-Dirac effect. It will be pointed out that the redistribution of momenta in the traveling wave can be considered as a transition from the resonant Kapitza-Dirac splitting to the Stern-Gerlach-type splitting. In Sec. IV will be discussed in details the temporal behavior of mean momenta corresponding to both internal energy levels. The results are summarized in Sec. V, where the possibility of experimental observation of this phenomenon is sketched too.

II. MOMENTUM DISTRIBUTIONS AND MEAN MOMENTA PER ATOMIC INTERNAL ENERGY LEVELS

Let us discuss the resonant interaction of a two-level atom with the radiation field [1]. For the sake of simplicity, suppose the field has a plane wavefront and linear polarization (these assumptions will be conserved for the standing wave, discussed in Sec. III, too). Let us suppose then that the field amplitude turns on instantly. Let the internal wave functions of a free two-level atom in ground (g) and excited (e) levels be $\varphi_g(\vec{\rho},t)$ and $\varphi_e(\vec{\rho},t)$, respectively, where $\vec{\rho}$ is the atomic internal coordinate (the radius vector of the optical electron, relative to the atomic center of mass). The wave function of an interacting atom will be [1]

$$\Psi = A\varphi_g(\vec{\rho},t) + B\varphi_e(\vec{\rho},t), \tag{1}$$

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where A and B are the probability amplitudes of the atom at the ground and excited internal energy levels correspondingly.

When taking into account the translational motion of the atomic center of mass, it is necessary to separate the corresponding parts (wave functions) in A and B coefficients. If, for example, the atom at the internal energy level has the well-defined value of momentum p, the corresponding wave function is given by the function

$$\chi(p)=\frac{1}{\sqrt{2\pi\hbar}}\exp\left(\frac{i}{\hbar}pz\right), \tag{2}$$

that is, by an exponential function with imaginary degree. In general, if an atom does not have definite values of momentum at any energy level, the coefficients A and B will be expressed by the series of $\chi(p)$ states:

$$A(t,z)=\int a(p,t)\chi(p)dp, \quad B(t,z)=\int b(p,t)\chi(p)dp, \tag{3}$$

with probability amplitudes $a(p,t)$ and $b(p,t)$ of the atom to have momentum p (at the moment t of time) and simultaneously to be at the ground or excited internal energy levels, correspondingly.

Inserting the expressions (1)–(3) into the quantum-mechanical determination of atom momentum

$$\langle p\rangle=\int \Psi^*\hat{p}\Psi d\vec{\rho}dz, \quad \int \Psi^*\Psi d\vec{\rho}dz=1, \tag{4}$$

and doing the standard transformations we arrive at

$$\langle p\rangle=\int |a(p,t)|^2pdp+\int |b(p,t)|^2pdp. \tag{5}$$

The first member specifies the contribution of translational states per ground internal energy level into the total momentum,

$$\langle p\rangle_g=\int |a(p,t)|^2pdp. \tag{6}$$

Accordingly, the second member specifies the excited level's state's contribution,

$$\langle p\rangle_e=\int |b(p,t)|^2pdp. \tag{7}$$

Both these momenta are time-dependent and their changes for interaction time t will be

$$\langle \Delta p\rangle_g=\int [|a(p,t)|^2-|a(p,0)|^2]pdp, \tag{8a}$$

$$\langle \Delta p\rangle_e=\int [|b(p,t)|^2-|b(p,0)|^2]pdp. \tag{8b}$$

When the atom interacts with the traveling wave, the internal ground level coefficient $a(p,t)$ relates with the excited internal level coefficient $b(p+\hbar k,t)$ (the spontaneous transitions are not taken into account). As a result we get a conserving quantity

$$|a(p,t)|^2+|b(p+\hbar k,t)|^2 = \text{const}=|a(p,0)|^2+|b(p+\hbar k,0)|^2 \tag{9}$$

[it can be checked by Eqs. (25)]. We can connect $\langle \Delta p\rangle_g$ to $\langle \Delta p\rangle_e$ by means of this relation:

$$\begin{aligned}\langle \Delta p\rangle_e&=\int [|b(p+\hbar k,t)|^2-|b(p+\hbar k,0)|^2] \times(p+\hbar k)d(p+\hbar k)\\ &=-\int [|a(p,t)|^2-|a(p,0)|^2](p+\hbar k)d(p+\hbar k)\\ &=-\langle \Delta p\rangle_g+\hbar k\int [|a(p,t)|^2-|a(p,0)|^2]dp\\ &=-\langle \Delta p\rangle_g+\hbar k\,\Delta n_g,\end{aligned} \tag{10}$$

where $\Delta n_g=-\Delta n_e=\int[|a(p,t)|^2-|a(p,0)|^2]dp =-\int[|b(p,t)|^2-|a(p,0)|^2]dp$ is the change of internal ground level's population, or which is the same, the population change Δn_e of the internal excited level with the opposite sign [see Eq. (19)]. From the equality of the first and last parts of Eq. (10), the well known inequality between the momenta of photon and total atom directly follows:

$$\langle \Delta p\rangle=\langle \Delta p\rangle_g+\langle \Delta p\rangle_e=\hbar k\,\Delta n_g\leqslant\hbar k. \tag{11}$$

Note, nevertheless, that this ''one photon demarcation'' pertains to the total momentum of the atom and not to the momentum per ground and excited internal levels, separately. Their changes, in accordance with Eqs. (8a) and (8b), in principle, can be arbitrary, depending on the distributions of $|a(p,t)|^2-|a(p,0)|^2$ and $|b(p,t)|^2-|b(p,0)|^2$ in the momentum space. From the expressions (8a) and (8b) also it is obvious that to get great values for $\langle \Delta p\rangle_g(\langle \Delta p\rangle_e)$, it is necessary that the distribution of $|a(p,t)|^2-|a(p,0)|^2$ [or $|b(p,t)|^2-|b(p,0)|^2$] be strictly nonsymmetric, relative to the replacement $p\rightarrow -p$, and to have a gathering in the range of great values of $|p|$.

And now let us show that the one-photon absorption/emission process in the field of traveling waves really allows a behavior, mentioned above. The Hamiltonian of the system, in the dipole approximation, can be written as

$$\hat{H}=\hat{H}_0-\hat{d}E(t,z), \tag{12}$$

where $\hat{H}_0$ is the free atom Hamiltonian, $\hat{d}$ is the dipole moment operator, and

$$\overline{E}(t,z)=\frac{\overline{E}}{2}\exp(ikz-i\omega t)+\text{c.c.}, \quad t>0 \tag{13}$$

is the electric field, whose frequency ω is equal to the Bohr transition frequency ω_0.

From the Schrödinger equation for $A(t,z)$ and $B(t,z)$ amplitudes we arrive at

$$i\frac{\partial A(t,z)}{\partial t} = -\nu \exp(-ikz)B(t,z), \tag{14a}$$

$$i\frac{\partial B(t,z)}{\partial t} = -\nu \exp(ikz)A(t,z), \tag{14b}$$

the Rabi solutions of which are [1]

$$A(z,t) = A(z,0)\cos \nu t + iB(z,0)\exp(-ikz)\sin \nu t, \tag{15a}$$

$$B(z,t) = B(z,0)\cos \nu t + iA(z,0)\exp(ikz)\sin \nu t, \tag{15b}$$

where $\nu = dE/2\hbar$ represents the Rabi frequency, $d = \langle \varphi_a | \hat{d} | \varphi_b \rangle$.

Performing the $\chi(p)$ expansion [see Eq. (3)] in Eqs. (15a) and (15b), we obtain

$$a(p,t) = a(p,0)\cos \nu t + ib(p+\hbar k,0)\sin \nu t, \tag{16a}$$

$$b(p,t) = b(p,0)\cos \nu t + ia(p-\hbar k,0)\sin \nu t. \tag{16b}$$

At first, it is readily verified that if the atom is at one of the energy levels before the interaction, the extraordinary things will not take place. Really, if for example $b(p,0)=0$, then

$$\langle \Delta p \rangle_g = (\cos^2 \nu t - 1)\int |a(p,0)|^2 p\,dp$$

$$= (\cos^2 \nu t - 1)\langle p \rangle_g|_{t=0}, \tag{17a}$$

$$\langle \Delta p \rangle_e = (1-\cos^2 \nu t)[\langle p \rangle_g|_{t=0} + \hbar k], \tag{17b}$$

that is, the contribution of momentum per each internal energy level evolves periodically and this is the evolution merely caused by a periodic exchange of the population between the internal energy levels [posed by the term $(1-\cos^2 \nu t)$]. Note also that in conditions under consideration the momentum distributions coincide with each other with a $\hbar k$ shift: $b(p+\hbar k,t) = ia(p,t)\tan \nu t$, as was mentioned in the Introduction.

The situation is totally diverse, if the atom is initially in a superposition state of the internal ground and excited levels; now the initial momentum distributions on the internal ground and excited levels in general are not required to be identical with the $\hbar k$ shift: $b(p,0) \neq \alpha\ a(p-\hbar k,0)$ (α is some constant, independent of p). Then, it unavoidably follows from Eqs. (16a) and (16b) that the optical transition, besides the changes in the internal energy levels' populations, also leads to periodic evolutions in the form of momentum distributions. Thus the atomic amplitudes $a(p,t)$ and $b(p,t)$ are not mutually proportional (with any constant shift).

To wash out the contributions, appropriate to evolution of the internal energy level populations, let us introduce a pair of new quantities, $\overline{p_g}$ and $\overline{p_e}$, which would be scaled in units of the level populations n_g and n_e, respectively:

$$\overline{p_g} = \langle p \rangle_g / n_g, \quad \overline{p_e} = \langle p \rangle_e / n_e, \tag{18}$$

$$n_g = \int |a(p,t)|^2 dp, \quad n_e = \int |b(p,t)|^2 dp. \tag{19}$$

Since these new quantities are already independent of the internal level populations, their possible evolutions would be stipulated by form deformations in the internal energy level momentum distributions. Afterwards we will call them mean momenta of an internal ground energy level ($\overline{p_g}$) and of an internal excited energy level ($\overline{p_e}$) correspondingly. Thereby, the total momentum of an atom, in addition to Eq. (5), can be represented in the more convenient form

$$\langle p \rangle = n_g \overline{p_e} + n_e \overline{p_g}. \tag{20}$$

These mean momenta, $\overline{p_g}$ and $\overline{p_e}$, remain constant, of course, if the atom is initially at one of the energy levels. They also remain constant if the initial distributions are mutually proportional with the constant $\hbar k$ shift:

$$b(p,0) = \alpha\, a(p-\hbar k,0). \tag{21}$$

Really, putting Eq. (21) in relations (16a) and (16b) and making the obvious substitutions, we arrive at

$$\overline{p_g} = \frac{|\cos \nu t - i\alpha \sin \nu t|^2 \int |a(p,0)|^2 p\,dp}{|\cos \nu t - i\alpha \sin \nu t|^2 \int |a(p,0)|^2 dp} = \overline{p_g}|_{t=0}$$

for the internal ground energy level, and $\overline{p_e} = \overline{p_e}|_{t=0}$ for the internal excited energy level. In these circumstances in Eq. (20) the time evolution exhibits only the internal energy level populations n_e and n_g.

In the general case, nevertheless, the state evolution is due to the interference of nonsimilarly distributed amplitude. The atomic amplitudes distributions per internal energy levels are not proportional to each other and subsequently the mean momenta $\overline{p_e}$ and $\overline{p_g}$ get temporal evolution, too.

For acquisition of more concrete and quantitative results, let us note that for the intentions of atom optics and interferometry [2], the coherent scattering of atoms in the resonant field of a standing wave is the routine for preparation of a large spreaded momentum distribution. The probability amplitudes, prepared in a such way, cannot satisfy the ''undesirable'' condition (21), in principle, since in the field of a standing wave, as is well known, any state with momentum p at one internal energy level is connected simultaneously to the two states at the other internal energy level with momenta $p-\hbar k$ and $p+\hbar k$. Therefore, any atom prepared by means of resonant Kapitza-Dirac effect, during its later interaction with the traveling wave, ought to implicitly change the momentum distributions at the internal energy levels with the results mentioned above.

III. THE CASE OF PREPARATION OF ATOMIC SUPERPOSITIONAL STATES BY SCATTERING IN THE FIELD OF RESONANT STANDING WAVES

Let us consider that the atom has a coherent interaction with the resonant ($\omega=\omega_0$) standing wave [3] during the time τ_s before the interaction with the traveling wave. We restrict ourselves to the relatively simple case when the interaction proceeds by the well known scheme of mutually orthogonal atom-standing wave beams. Moreover, the Raman-Nath approximation will be applied, which permits us to put out the kinetic energy term in the Hamiltonian from the problem at hand [note that the kinetic energy term has not been included in Eq. (12) either]. Although the scheme of calculations is well known and is represented in detail see, for example, [2,3]) we find it convenient to give here an account of the main intermediate formulas, too.

To describe the interaction in the preparing standing wave, the electric field (13) in the Hamiltonian (12) must be exchanged by

$$E(t,z)=E_s\cos kz\exp(-i\omega t)+\text{c.c.},\quad -\tau_s\leqslant t\leqslant 0. \tag{22}$$

As a consequence, the atomic amplitudes $A_s(z,t)$ and $B_s(z,t)$ have to satisfy Eqs. (14a) and (14b) type equations where the following replacements must be performed: $\nu\rightarrow 2\nu_s=2dE_s/\hbar$ (which is the mean Rabi frequency in the standing wave), $\exp(\pm ikz)\rightarrow\cos kz$. Allowing that the atom has been at the ground level before the interaction ($t<-\tau_s$), we arrive at

$$A_s(z,t)=\cos[2\nu_s(t+\tau_s)\cos kz], \tag{23a}$$

$$B_s(z,t)=i\sin[2\nu_s(t+\tau_s)\cos kz]. \tag{23b}$$

These amplitudes at the moment $t=0$, when the standing wave is turned off, just represent the initial amplitudes $A(z,0)$ and $B(z,0)$ for an interaction with the traveling wave [see the formulas (15a) and (15b)]. Making their $\chi(p)$ expansions [4],

$$A_s(z,0)=\cos(2\nu_s\tau_s\cos kz)=\sum_{m=-\infty}^{\infty}i^{2m}J_{2m}(2\nu_s\tau_s)\exp(i2mkz), \tag{24a}$$

$$B_s(z,0)=i\sin(2\nu_s\tau_s\cos kz)=\sum_{m=-\infty}^{\infty}i^{2m+1}J_{2m+1}(2\nu_s\tau_s)\exp(i(2m+1)kz), \tag{24b}$$

where m is the number of photons reemited from one into the other of the counterpropagating waves, $J_m(x)$ is a Bessel function. For the atomic center-of-mass motion probability amplitudes $a(p,t)$ and $b(p,t)$, we get the following expressions:

$$a(2m\hbar k,t)=i^{2m}[\cos\nu t\,J_{2m}(2\nu_s\tau_s)-\sin\nu t\,J_{2m+1}(2\nu_s\tau_s)], \tag{25a}$$

$$b((2m+1)\hbar k,t)=i^{2m+1}[\cos\nu t\,J_{2m+1}(2\nu_s\tau_s)+\sin\nu t\,J_{2m}(2\nu_s\tau_s)], \tag{25b}$$

$$a((2m+1)\hbar k,t)=b(2m\hbar k,t)=0. \tag{25c}$$

We see that the superposition state, created as a result of the interaction with the standing wave, represents discrete manifolds of states, where the space between the adjacent values of momentum is $2\hbar k$, herewith the manifolds for the ground and excited internal energy levels are totally shifted with respect to each other by $\hbar k$ (half of $2\hbar k$) [3].

The formulas (25a)–(25c) contain explicitly the sought-after result about the evolution of momentum distributions. To exhibit this evolution, let us first note that the initial momentum distribution for both internal energy levels is symmetric relative to value $p=0$. Really, they are specified by $i^{2m}J_{2m}(\cdot)$ and $i^{2m+1}J_{2m+1}(\cdot)$ functions for the ground and excited energy levels, respectively, and are symmetric relative to $2m\rightarrow -2m$, $2m+1\rightarrow -(2m+1)$ transformations, that is, relative to the value $m=0$ ($p=0$). This symmetry signifies that the momentum per each internal energy level [both for Eq. (5) and Eq. (20) values] is zero [3] before the interaction with the traveling wave. Nevertheless, in accordance to Eqs. (25a) and (25b), the symmetry breaks under the ''influence'' of the traveling wave: the one-photon absorption/emission process gives the beginning of asymmetric transformations in the form of momentum distributions, periodically running in opposite directions for the ground and excited internal energy levels.

A typical form of the initial distributions and the following redistributions (due to a single-photon process) are depicted in Figs. 1(a) and 1(b) for the ground and excited energy levels by turns. Single-photon large-scale changes are apparent.

Now let us notice that in conditions of Figs. 1(a) and 1(b) we get almost one-side distributions: the translational states with $n>0$ for the ground internal energy level only and the translational states with $n<0$ for excited internal level only. So, the state of the total atom has been split into two subgroups, where one subgroup represents the ground-level atoms with positive values of the momentum, and the second subgroup represents the opposite, namely the excited-level atoms with negative values of momentum. Of course, this is a Stern-Gerlach-type splitting. That is, a one-photon optical transition implements the resonant Kapitza-Dirac splitting into the Stern-Gerlach-type splitting.

The phenomenon of one-photon coherent accumulation of the momentum on the internal energy levels (OP-CAMEL) can get some expansion, if the initial momentum distributions will be taken in an asymmetric form. This kind of distribution can also be obtained by the standing wave, but only if a traveling wave is preceding it [5]. Such a sequence of pulses is obtained if the standing wave is formed by means of the reflection of a laser pulse from the mirror (see, for example, [6]). To avoid the overloading of the text we do not

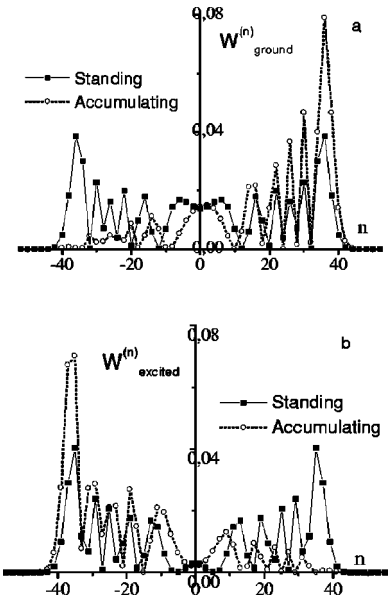

FIG. 1. Probability distribution of definite-momentum states. (Upper plot) $W^{(m)}_{\text{ground}}=|a(2m\hbar k,t)|^2$ for ground internal energy level and (lower plot) $W^{(m)}_{\text{excited}}=|b((2m+1)\hbar k,t)|^2$ for excited internal energy level, prepared symmetrically in the momentum space by the interaction with the standing wave. The chosen parameters are $2\nu_s\tau_s=40$, consequently $|A(-\tau_s)/\chi(P_0)|^2=1$, $|B(-\tau_s)/\chi(P_0)|^2=0$, $\nu t=\pi/4$.

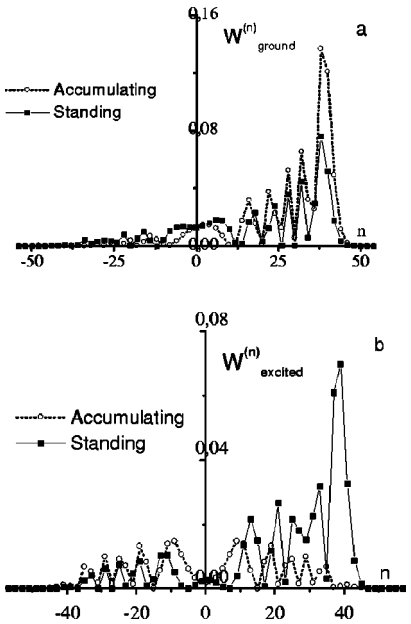

FIG. 2. Probability distribution of definite-momentum states. (Upper plot) $W^{(m)}_{\text{ground}}$ for internal ground energy level and (lower plot) $W^{(m)}_{\text{excited}}$ for internal excited energy level, prepared asymmetrically by the interaction with the standing wave. The chosen parameters are $2\nu_s\tau_s=40$, consequently $|A(-\tau_s)/\chi(P_0)|^2=1/2$, $|B(-\tau_s)/\chi(P_0)|^2=1/2$, $\nu t=\pi/4$.

give the formulas, and the behavior of the OP-CAMEL in these conditions will be given only by some graphs. In Figs. 2(a) and 2(b), the case is presented when the momentum distribution in resonant Kapitza-Dirac splitting is maximum asymmetric. As is seen from the graphs, in this special case OP-CAMEL appears already as an accumulation of an asymmetry for one (ground) internal energy level at the appearance of its suppressing for the other (excited) internal energy level.

IV. TIME EVOLUTION OF MEAN MOMENTUM PER GROUND AND EXCITED INTERNAL ENERGY LEVELS IN THE FIELD OF TRAVELING WAVES

Let us now take the case of preparation by the standing wave and discuss the evolution of momenta $\overline{p_g}$ and $\overline{p_e}$. By means of expressions for the quantities defining $\overline{p_g}$ and $\overline{p_e}$ [see Eqs. (18), (19), (6), and (7)] we will have

$$\langle p\rangle_g=\hbar k\sum_{m=-\infty}^{\infty}2m[\cos\nu t\,J_{2m}(u)-\sin\nu t\,J_{2m+1}(u)]^2$$
$$=-\hbar k\left[\frac{1-J_0(2u)}{2}\sin^2\nu t+\frac{u-J_1(2u)}{4}\sin 2\nu t\right], \tag{26}$$

$$n_g=\sum_{m=-\infty}^{\infty}[\cos\nu t\,J_{2m}(u)-\sin\nu t\,J_{2m+1}(u)]^2$$
$$=\frac{1}{2}+\frac{J_0(2u)}{2}\cos 2\nu t-\frac{J_1(2u)}{2}\sin 2\nu t \tag{27}$$

for the ground internal energy level, and

$$\langle p\rangle_e=\hbar k\sum_{m=-\infty}^{\infty}(2m+1)[\cos\nu t\,J_{2m+1}(u)$$
$$+\sin\nu t\,J_{2m+1}(u)]^2$$
$$=\hbar k\left[\frac{1+J_0(2u)}{2}\sin^2\nu t+\frac{u+J_1(2u)}{4}\sin 2\nu t\right], \tag{28}$$

$$n_e=\sum_{m=-\infty}^{\infty}[\cos\nu t\,J_{2m+1}(u)+\sin\nu t\,J_{2m}(u)]^2$$
$$=\frac{1}{2}-\frac{J_0(2u)}{2}\cos 2\nu t+\frac{J_1(2u)}{2}\sin 2\nu t=1-n_g \tag{29}$$

for the excited internal energy level. Here $u=2\nu_s\tau_s$. The last forms of Eqs. (26)–(29) are obtained by using the formulas of summations of Bessel functions [7] just as $\langle p\rangle_g|_{t=0}=0$, $\langle p\rangle_e|_{t=0}=0$ (the same for $\overline{p_g}$ and $\overline{p_e}$), so their values at any next moment t present simultaneously their changes for the time t: $\langle\Delta p\rangle_g=\langle p\rangle_g$, $\langle\Delta p\rangle_e=\langle p\rangle_e$.

In Figs. 3(a) and 3(b), the temporal evolutions of momenta are given, corresponding to each internal energy level, while they are interacting with traveling (accumulating) waves. Population changes, which are also responsible for the momentum time evolutions in general, are depicted in the figure by dashed lines. In the represented case during the interaction with the traveling wave, the populations are practically constant, which follows from Eqs. (27) and (29), too,

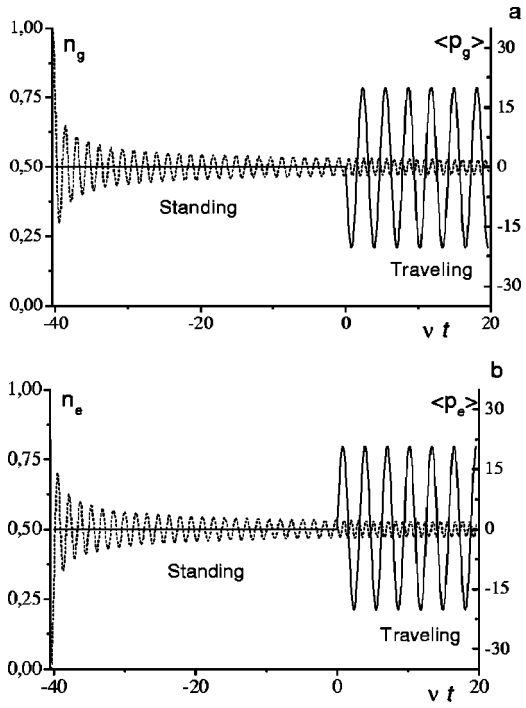

FIG. 3. Temporal behavior of momentum per (upper plot) ground internal energy level and (lower plot) excited internal energy level. Time interval includes all parts of interaction: with preparing-standing and accumulating-traveling waves. All parameters have the same values as in Fig. 1.

if we take into account that $J_{0,1}(2u) \ll 1$ for $u \gg 1$. Respectively, the temporal evolutions of the mean momenta $\overline{p_g}$ and $\overline{p_e}$ are due to the redistributions of momentum at the energy levels only and are shown in Figs. 4(a) and 4(b). The parameters of the preparing standing wave are the same as in Figs. 1(a) and 1(b), where the distance between the left-hand and

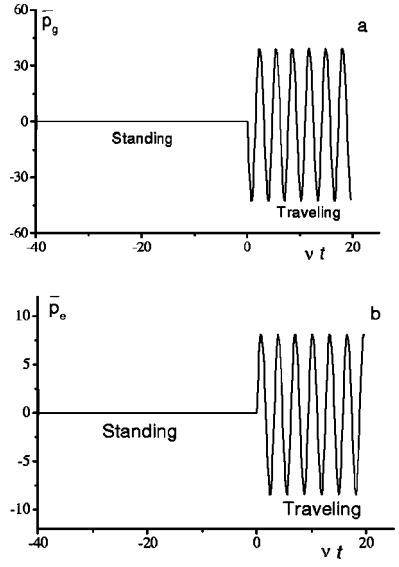

FIG. 4. Temporal behavior of mean momentum per (upper plot) ground internal energy level and (lower plot) excited internal energy level.

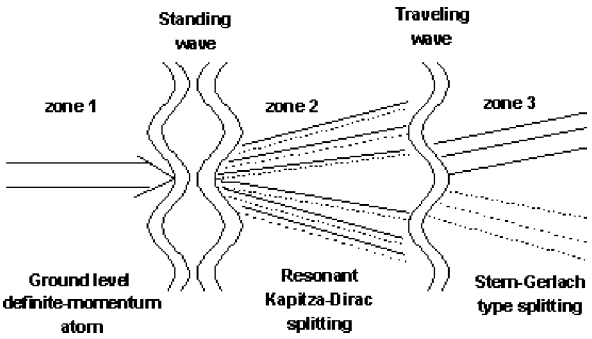

FIG. 5. The ground internal energy level; definite momentum state of an atom (zone 1) transforms into a superposition one by coherent interaction with a resonant standing wave (zone 2). The next interaction with the traveling wave leads to large-scale changes in atomic momentum distributions per internal energy level. Solid lines present the ground-level and the dotted lines present the excited-level atomic states.

right-hand maxima (the width of momentum distribution) is about 70 $\hbar k$. Such magnitudes for the resonant Kapitza-Dirac splitting are totally in the limits of experimental realizations [8].

Note also that the comparison of the deviation of $\overline{p_g}$ or $\overline{p_e}$ [from Figs. 4(a) and 4(b)] and the width of momentum distribution [from Figs. 1(a) and 1(b)] show the same order of magnitude for them. Since the width of momentum distribution has multiphoton nature (created by means of a multiphoton process of reemission of photons from one wave into the counterpropagating one), the large-scale variations in OP-CAMEL may be called as ''multiphoton.''

Multiphoton OP-CAMEL manifests itself in Eqs. (26) and (28) in the following way. When the initial momentum distribution is sufficiently widespread, that is, $\Delta p \gg \hbar k$, then $u = 2\nu_s \tau_s \gg 1$ (because in the theory of the resonant Kapitza-Dirac effect the connection between momentum width δp and the number of Rabi-flops $2\nu_s \tau_s$ is $\delta p \approx 2\nu_s \tau_s \hbar k$). Taking also into account that $J_{1,0}(x) \leqslant 1$, we obtain that the members $\frac{1}{4}\hbar k u \sin 2\nu t$ in Eq. (26) and $-\frac{1}{4}\hbar k u \sin 2\nu t$ in Eq. (28) stand out as the prevailing terms for $\langle p \rangle_g$ and $\langle p \rangle_e$, respectively:

$$\langle p \rangle_g \approx -\frac{u}{4}\hbar k \sin 2\nu t, \quad \langle p \rangle_e \approx \frac{u}{4}\hbar k \sin 2\nu t.$$

Since $u \gg 1$, we see that the changes of each momentum per Rabi period, being $(u/2)\hbar k$, exceed greatly the photon's momentum $\hbar k$.

V. SUMMARY

A simple theoretical consideration of the optical transition for general conditions, when the atom in the superposition state of ground and excited internal energy levels initially has different momentum distributions, shows that the one-photon optical transition leads to radical asymmetric changes in momentum distributions at each internal energy level. In other words, a photon's change of the mean momentum of

each internal energy level is more than the photon's own momentum.

For an important case, when the preliminary superposition state of the atom is prepared by coherent scattering at the resonant standing wave, the phenomenon for definite intervals of time can be presented as a transition from resonant Kapitza-Dirac splitting of atomic translational states into Stern-Gerlach-type splitting. This is schematically depicted in Fig. 5.

Finally, let us make some remarks on the possibility of experimental observation of the phenomenon. First, let us notice that the ''nonoptical'' methods, which detect the total atom (for example, the ''hot-wire'' method), cannot be used for this purpose, because the phenomenon deals with each individual internal energy level; the momentum distribution of the total atom does not change, or rather, it changes only in one-photon momentum limits.

It is preferable to use detecting methods, which deal only with one of the resonantly connected internal energy levels, such as the adjacent optical transitions. Then the phenomenon will appear as a pronounced asymmetry in the profile of Doppler broadening, relative to Bohr frequency. Another possibility we see in using the long-living energy levels is that thus far the atomic translational states can be distinguished in space before the spontaneous emission (zone 3 in Fig. 5). In this case the space-sensitive schemes of spontaneous-emission collection or probe pulse absorption will result in a desirable outcome.

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