Large Scale (Multiphoton) Evolution of Atomic Energy Level’s Mean Momentum for Single-Photon Absorption/Emission Process

A. Zh. Muradyan\(^1\), H. L. Haroutyunyan\(^2\)

\(^1\)Department of Physics, Yerevan State University, I Alex Manukian, Yerevan Armenia

\(^2\)Engineering Center of Armenian National Academy of Sciences, Ashtarak-2, 378410 Armenia;

E-mail: muradyan@ec.sci.am; yndanfiz@ysu.am.

It is shown that a two-level atom, being initially in general superposition state of ground and excited energy levels with mutually different momentum distributions there, gets a large scale evolution in the energy levels momentum distribution. As a consequence the mean momentum of each individual energy level also gets large scale changes, more than the own momentum \(\hbar k\) of the photon. Thoroughly is discussed the special case, when 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 that in such conditions the mentioned phenomenon can be presented as a transformation of the resonant Kapitza-Dirac splitting of atomic states into the Stern-Gerlach type splitting.

I. INTRODUCTION

Interaction of a two-level system with a plane traveling wave leads to one-photon transitions between energy levels and consequent changes of total atomic momentum in limits of one photon momentum \(\hbar k\). What can be said about distributions and mean values of momentum for each individual atomic level? The answer is well-known and is trivial, if the atom before the interaction is on one of the energy levels: distribution on the other level sets shift \(\hbar k\) and mean values of momentum distinguished also by \(\hbar k\); \(\overrightarrow{p}_g = \overrightarrow{p}_g + \hbar k\), where \(\overrightarrow{p}_g\) and \(\overrightarrow{p}_e\) are mean values of momentum on ground and excited levels consequently (1D case). What would we have in general case, that is, when the atom before the interaction with the travelling wave is in the superposition state of ground and excited levels with mutually different momentum distributions there? In the following just this question will be elucidated in details: in general form in Sec. 2, and the special case of preparing standing wave, latter.

In Sec. 2 it will be shown that in one-photon absorption/emission process, in general, large scale redistributions of the energy level’s momenta takes place. This entails the time-changes of the \(\overrightarrow{p}_g\), \(\overrightarrow{p}_e\) and these changes can greatly exceed the value of one photon momentum. In general, the relations \(\overrightarrow{p}_e = \overrightarrow{p}_g + \hbar k\) are not true, too.

In Sec. 3 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 resonant standing wave, which is being often referred to us the resonant Kapitza-Dirac effect. Will be pointed out, that the redistribution of momenta in the travelling can be received as a transition from the resonant Kapitza-Dirac splitting to the Stern-Gerlach type splitting. In Sec. 4 will be discussed the temporal behavior of the energy level’s mean momenta. The results are summarized in Sec. 5, where the possibility of experimental observation of this phenomenon will be sketched too.

II. DISTRIBUTIONS AND MEAN VALUES OF MOMENTUM ON THE GROUND AND EXCITED ENERGY LEVELS IN THE ABSORPTION/EMISSION OF ONE PHOTON

Let’s discuss the resonant interaction of a two-level atom with the radiation field \(\|\). For the sake of simplicity, suppose the field has plane wavefront, linear polarization (these assumptions will be conserved for the standing wave, discussed in Sec.3), turns on instantaneously and after on its amplitude remains constant. Let the wave functions of free two-level atom’s ground \((g)\) and excited \((e)\) levels be \(\varphi_g(\overrightarrow{\varphi}, t)\) and \(\varphi_e(\overrightarrow{\varphi}, t)\) respectively, where \(\overrightarrow{\varphi}\) is atomic internal coordinate (radius-vector of the optical electron relative to the atomic center-of-mass). The wave function of interacting atom will be \(\|\).

\(^1\)Because of not completeness of the literature under the hand, we are not sure that this simple phenomenon isn’t known in quantum theory of atom-photon interactions or, particularly, in the interferometry of atomic matter waves. One of main purposes of this article is our request to the leading specialists to respond and if the phenomenon is known, send the corresponding references.
\[
\Psi = A \varphi_g(\overrightarrow{\rho}, t) + B \varphi_e(\overrightarrow{\rho}, t)
\]

(1)

where \(A\) and \(B\) are the probability amplitudes of the atom to be on the ground and excited levels correspondingly.

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

\[
\chi(p) = \frac{1}{\sqrt{2\pi\hbar}} \exp\left(\frac{i}{\hbar} p z\right),
\]

(2)

that is, by the exponential function with imaginary degree. In general, when on energy levels an atom hasn’t definite values of momentum, \(A\) and \(B\) coefficients will be expressed by the series of \(\chi(p)\)-states:

\[
A(t, z) = \int a(p, t)\chi(p)dp, \quad B(t, z) = \int b(p, t)\chi(p)dp,
\]

(3)

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

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

\[
\langle p \rangle = \int \Psi^* \hat{p} \Psi d\overrightarrow{\rho} dz,
\]

\[
\int \Psi^* \Psi d\overrightarrow{\rho} dz = 1
\]

(4)

and doing the standard transformations we arrive to

\[
\langle p \rangle = \int |a(p, t)|^2 p dp + \int |b(p, t)|^2 p dp.
\]

(5)

The first member specifies the contribution of the ground energy level in total momentum,

\[
\langle p \rangle_g = \int |a(p, t)|^2 p dp.
\]

(6)

Accordingly, the second member specifies the excited level’s contribution,

\[
\langle p \rangle_e = \int |a(p, t)|^2 p dp
\]

(7)

Both these momenta are time-dependent and their change for \(t\) time of interaction will be

\[
\langle \Delta p \rangle_g = \int \left(|a(p, t)|^2 - |a(p, 0)|^2\right) p dp,
\]

\[
\langle \Delta p \rangle_e = \int \left(|b(p, t)|^2 - |b(p, 0)|^2\right) p dp.
\]

(8a)

(8b)

When the atom interacts with the travelling wave, the coefficient \(a(p, t)\) of ground level relates with the coefficient \(b(p + \hbar k, t)\) of excited level (spontaneous transitions are not taken into account) (Fig. 1). As a result we get

\[
|a(p, t)|^2 + |b(p + \hbar k, t)|^2 = \text{const} = |a(p, 0)|^2 + |b(p + \hbar k, 0)|^2.
\]

(9)

By this relation we can connect \(\langle \Delta p \rangle_g\) with \(\langle \Delta p \rangle_e\):

\[
\langle \Delta p \rangle_e = \int \left(|b(p + \hbar k, t)|^2 - |b(p + \hbar k, 0)|^2\right) (p + \hbar k) d(p + \hbar k) =
\]

\[
= - \int \left(|a(p, t)|^2 - |a(p, 0)|^2\right) (p + \hbar k) d(p + \hbar k) =
\]

\[
= - \langle \Delta p \rangle_g + \hbar k \int \left(|a(p, t)|^2 - |a(p, 0)|^2\right) dp =
\]

\[
= - \langle \Delta p \rangle_g + \hbar k \Delta n_g,
\]

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

$$\langle \Delta p \rangle = \langle \Delta p \rangle_g + \langle \Delta p \rangle_e = \hbar k \Delta n_g \leq \hbar k ;$$  \hfill (11)

Let, nevertheless, note that this $<\text{one photon demarcafation}>>$ pertains to the total momentum of the atom, but not to the momentum of ground and excited levels, separately. Their changes, in accordance with [14] and [31], in principle, may be arbitrary, depending on distributions of $|a(p,t)|^2 - |a(p,0)|^2$ and $|b(p,t)|^2 - |b(p,0)|^2$ in momentum space. From the expressions (14a) and (14b) also is obvious that for $\langle \Delta p \rangle_g$ (or $\langle \Delta p \rangle_e$) to get great values, is necessary the distribution of $|a(p,t)|^2 - |a(p,0)|^2$ (of $|b(p,t)|^2 - |b(p,0)|^2$) to be strictly non-symmetric, relative to the replacement $p \rightarrow -p$ and to have a gathering in the range of great values of $|p|$.

And now let’s show that one photon absorption/ emission process in the field of travelling wave really allows a behavior, mentioned above. The Hamiltonian of the system, in dipole approximation, may be written as

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

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

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

is the electric field, whose $\omega$ frequency is equal to the $\omega_0$ frequency of Bohr transition.

From Schrodinger equation for $A(t, z)$ and $B(t, z)$ amplitudes we arrive to

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

$$i \frac{\partial B(t, z)}{\partial t} = -\nu \exp(ikz)A(t, z),$$  \hfill (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,$  \hfill (15a)
$$B(z, t) = B(z, 0) \cos \nu t + iA(z, 0) \exp(ikz) \sin \nu t,$$  \hfill (15b)

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

Performing $\chi(P)$-expansion (see (3)) in (15a) and (15b), we obtain

$$a(p, t) = a(p, 0) \cos \nu t + ib(p + \hbar k, 0) \sin \nu t,$$  \hfill (16a)
$$b(p, t) = b(p, 0) \cos \nu t + ia(p - \hbar k, 0) \sin \nu t,$$  \hfill (16b)

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

$$\langle \Delta p \rangle_g = \langle \cos^2 \nu t - 1 \rangle \int |a(p,0)|^2 pdp = (\cos^2 \nu t - 1) \langle P \rangle_g \big|_{t=0}$$  \hfill (17a)
$$\langle \Delta p \rangle_e = (1 - \cos^2 \nu t) \left[ \langle P \rangle_g \big|_{t=0} + \hbar k \right],$$  \hfill (17b)

that is the contribution of momentum per energy level evolves periodically and this is the evolution merely caused by periodic exchange of population between the 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 shift $\hbar k$: $b(p, t) = i a(p, t)$ $\text{tg} \nu t$, as was mentioned in Introduction.

The situation is totally diverse, if the atom is initially in superposition state of ground and excited levels, because now the initial momentum distributions on the ground and excited levels are not under necessity to be identical with shift $\hbar k$: $b(p, 0) \neq \alpha a(p - \hbar k, 0)$ in general ($\alpha$ is some constant, independent of $p$). Then, it follows unavoidably from (16a) and (16b) that the optical transition, besides the changes on the energy levels’populations, leads also to
III. PREPARATION OF SUPERPOSITIONAL STATES ON ATOMIC GROUND AND EXCITED LEVELS BY MEANS OF SCATTERING IN THE FIELD OF RESONANT STANDING WAVE

Let before the interaction with the travelling wave, during time $\tau_s$, the atom had coherent interaction with the resonant ($\omega = \omega_0$) standing wave $|\alpha\rangle$. 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 employed, which permit to put out from the problem at hand the kinetic energy term in the Hamiltonian (note that the kinetic energy term has not been included into (12) too). Although the scheme of calculation is well known and presented in details (see, for example, in [2], [3], we find it convenient to give an account of main intermediate formulas, too.

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

$$E(t, z) = E_s \cos k z \exp(-i \omega t) + c.c., \quad -\tau_s \leq t \leq 0. \quad (22)$$

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

\[ A_s(z, t) = \cos(2\nu s(t + \tau) \cos kz), \quad B_s(z, t) = i \sin(2\nu s(t + \tau) \cos kz). \]  

These amplitudes at the moment \(t = 0\), when the standing wave is turned off, just present the initial amplitudes \(A(z, 0)\) and \(B(z, 0)\) for interaction with the travelling wave (see the formulas \([15a, 15b]\)). Using their \(\chi(p)\)-expantions \([4]\).

\[ A_s(z, 0) = \cos(2\nu s \tau_s \cos kz) = \sum_{m = -\infty}^{\infty} i^{2m} J_{2m}(2\nu s \tau_s) \exp(i2mkz), \]  
\[ B_s(z, 0) = i \sin(2\nu s \tau_s \cos kz) = \sum_{m = -\infty}^{\infty} i^{2m+1} J_{2m+1}(2\nu s \tau_s) \exp(i(2m+1)kz), \]

where \(J_n(x)\) is Bessel function, for atomic center-of-mass motion probability amplitudes \(a(p, t)\) and \(b(p, t)\) we get the following expressions:

\[ a(2m\hbar k, t) = i^{2m} \left[ \cos \nu t \ J_{2m}(2\nu s \tau_s) - \sin \nu t \ J_{2m+1}(2\nu s \tau_s) \right], \]  
\[ b((2m + 1)\hbar k, t) = i^{2m+1} \left[ \cos \nu t \ J_{2m+1}(2\nu s \tau_s) + \sin \nu t \ J_{2m+1}(2\nu s \tau_s) \right], \]  
\[ a((2m + 1)\hbar k, t) = b(2m\hbar k, t) = 0 \]

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

The formulas \([25a]-[25c]\) contain explicitly the seeking result about the evolution of momentum distributions. To exhibit this evolution, let first note that the initial momentum distribution for both energy levels are symmetric relative to value \(0\). Really, they are specified by \(i^{2m} J_{2m}(\nu t \ J_{2m+1}(2\nu s \tau_s))\) functions for ground and excited energy levels respectively and are symmetric, relative to \(2m \to -2m\), \(2m+1 \to -(2m+1)\) transformations, that is just relative to value \(m = 0\) \((p = 0)\). This symmetry signifies that the momentum of each energy level (as the incremental \([3]\), as the mean \([20]\) values) is zero \(\hbar k\) before the interaction with the travelling wave. Nevertheless the symmetry breaks under the influence of travelling wave: one photon absorption/emission process, in accordance with \([25a]\) and \([25b]\), gives the beginning of asymmetric transformations in the form of momentum distributions, periodically running in opposite directions for ground and excited energy levels.

A typical form of initial distributions and the following redistributions (due to single-photon process) are depicted on Fig. 2 and Fig. 3 for ground and excited energy levels consequently. Single-photon large-scale changes are apparent.

Now let notice that in conditions of Fig. 2 and Fig. 3 we get almost one-side distributions: the translational states with \(n > 0\) for ground energy level only, and the translational states with \(n < 0\) for excited energy level only. So, the state of total atom has been exploited into two sub-groups, where one sub-group presents ground-level atoms with negative values of momentum, and the second sub-group presents vice-versa, the excited-level atoms with negative values of momentum. Of course, this is a Stern-Gerlach type splitting. That is, one-photon optical transition implements the resonant Kapitza-Dirac splitting into the Stern-Gerlach type splitting.

The phenomenon of one-photon coherent accumulation of momentum on energy levels (OP-CAMEL), may get some expansion in exhibition, if the initial momentum distributions would be taken in asymmetric form. This kind of distributions also can be obtained by the standing wave, but if some travelling wave is previous to it \([4]\). Such sequence of pulses is got, if the standing wave is formed by means of reflection of a laser pulse from the distanted (from atomic beam) mirror (see, for example \([1]\)). To avoid the overloading of the text we aren’t going to bring formulas and the behavior of the OP-CAMEL in these conditions will be given only by some graphs. In Fig. 4 and Fig. 5 is presented the case 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 a accumulation of asymmetry on one (ground) energy level for account of its suppressing on the other (excited) energy level.

**IV. TIME EVOLUTION OF MEAN MOMENTUM ON GROUND AND EXCITED ENERGY LEVELS IN THE FIELD OF TRAVELLING WAVE**

Let us now return to the preparation only 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 \([13]\) and \([4]\)) we will have
\[ \langle p \rangle_g = \hbar k \sum_{m = -\infty}^{\infty} 2m [\cos \nu t J_{2m}(u) - \sin \nu t J_{2m+1}(u)]^2 = \]
\[ = -\hbar k \left[ \frac{1 - J_0(2u)}{2} \sin^2 \nu t + \frac{u - J_1(2u)}{4} \sin 2\nu t \right] \]

\[ n_g = \sum_{m = -\infty}^{\infty} [\cos \nu t J_{2m}(u) - \sin \nu t J_{2m+1}(u)]^2 = \]
\[ = \frac{1}{2} + \frac{J_0(2u)}{2} \cos 2\nu t - \frac{J_1(2u)}{2} \sin 2\nu t \]
on the ground energy level, and

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

\[ n_e = \sum_{m = -\infty}^{\infty} [\cos \nu t J_{2m+1}(u) + \sin \nu t J_{2m}(u)]^2 = \]
\[ = \frac{1}{2} - \frac{J_0(2u)}{2} \cos 2\nu t + \frac{J_1(2u)}{2} \sin 2\nu t = 1 - n_g \]
on the excited energy level. Here \( u = 2\nu_s \tau_s \). The last forms of (27)-(29) are got by using for formulas of summations of Bessel functions \( J_0, J_1 \). As \( \langle p \rangle_g |_{t=0} = 0 \), \( \langle p \rangle_e |_{t=0} = 0 \), (the same for \( n_g \) and \( n_e \)), then their values at any next moment \( t \) present simultaneously their changes for the time \( t \). \( \langle \Delta p \rangle_g = \langle p \rangle_g - \langle p \rangle_e = \langle p \rangle_e \). In Fig. 6 and 7 are given the temporal evolutions of momenta, corresponding to each energy level, while they are interacting with travelling (accumulating) wave. Population changes, which also are responsible for the evolutions in general, are depicted in figures by dashed lines. In presented case the populations are constant in practice, which follows from (27) and (29) too, if we take into account \( J_{0,1}(2u) << 1 \) for \( u >> 1 \). And, finally, the temporal evolution of mean momentum \( \overline{p}_g \) and \( \overline{p}_e \) conditioned solely by redistributions of momentum the energy levels, is shown already in Fig. 8 and 9. Parameters of the preparing standing wave are the same as in Fig. 2, where the distance between left-hand and right-hand maxima (the width of momentum distribution) is about \( 70 \hbar k \). Such magnitudes for the resonant Kapitza-Dirac splitting are totally in limits of experimental realizations \( \{\} \). Note also, that the comparison of the deviation of \( \overline{p}_g \) or \( \overline{p}_e \) (from the Fig. 8 and 9) and the width of momentum distribution (from the \( \{\} \)) shows the same order of magnitude for them. Since the width of momentum distribution has multiphoton nature (created by means of multiphoton process of reemission of photons from one wave into the counterpropagating one), the large-scale variations in OP-CAMEL may be called as \( \langle \langle \text{multiphoton} \rangle \rangle \).

Multiphoton OP-CAMEL manifests itself in (26), (28) formulas in following way. When the initial momentum distribution is sufficiently widespread, that is \( \Delta p \gg \hbar k \), then \( 2\nu_s \tau_s \gg 1 \) (from the theory of resonant Kapitza-Dirac effect the connection between the momentum width \( \Delta P \) and the number of Rabi-flops \( 2\nu_s \tau_s \approx \Delta p \approx 2\nu_s \tau_s \hbar k \)). Under this condition the members \(-\frac{1}{4} \hbar k u \sin \nu t \) in (26) and \(-\frac{1}{4} \hbar k u \sin \nu t \) in (28), being proportional to \( 2\nu_s \tau_s >> 1 \), stand out as the previal terms and just present the multiphoton OP-CAMEL.

V. SUMMARY

A simple theoretical consideration of optical transition in general conditions, when the atom initially has in superposition state of ground and excited energy levels different momentum distributions, shows that one one-photon optical transition leads to radical asymmetric changes in momentum distributions on each energy level. Saying in images, in general one photon changes the mean momentum of energy level more than photon’s own momentum.

In important case, when the preliminary superposition state of the atom is being prepared by coherent scattering at the resonant standing wave, the phenomenon can be presented as a transition from resonant Kapitza-Dirac splitting of atomic translational states to Stern-Gerlach type splitting. This is schematically depicted in Fig. 10.
Finally, let’s give an account of some remarks about the possibility of experimental observation of phenomenon. First let’s notice that the <<non-optical>> methods, which registrate the total atom (for example <<hot-wire>> method), can’t be used for this purpose, because the phenomenon deals with each individual energy level; the momentum distribution of total atom doesn’t change or which will be more right, it changes only in one-photon momentum limits.

It is preferable to use registration methods, that will deal only with one of resonantly connected energy levels, such as the adjacent optical transitions. Then the phenomenon will appear itself as a pronounced asymmetry in profile of Doppler broadening, relative to Bohr frequency. Another possibility we see in using of long-living energy levels, so that the atomic translational states can be distinguished in space till the spontaneous emission (C zone in Fig. 10). In this case the space-sensitive schemes of spontaneous emission collection or probe pulse absorption will bring to desirable result.

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[6] See Ref. 4, p.p. 665-667.
[7] See Ref. 4, p.p. 665-667.
[8] See, for example, last ref. in [3].

FIG. 1. The scheme of interaction of a two-level atom with the resonant travelling wave. The initial state of the atom is prepared by the standing wave and presents a manifold of discrete translational states per each energy level.

FIG. 2. Probability distribution $W^{(n)}_{\text{ground}}(\nu)$ of definite-momentum states on ground level prepared by the interaction with the standing wave. Integer $n$ determines the value of momentum ($p = n\hbar k$). Momentum distribution before the travelling wave was symmetric. The chosen parameters are $2\nu\tau_s = 40$, consequently $\left| \frac{A(-\tau_s)}{A(\tau_s)} \right|^2 = 1$, $\left| \frac{B(-\tau_s)}{B(\tau_s)} \right|^2 = 0$, $\nu t = 20$.

FIG. 3. Probability distribution $W^{(n)}_{\text{excited}}(\nu_t)$ of definite-momentum states in excited level after the interaction with the standing and travelling waves. Integer $n$ determines the value of momentum ($p = n\hbar k$). Momentum distribution before the travelling wave was symmetric. The chosen parameters are $\nu_t\tau_s = 20$, consequently $\left| \frac{A(-\tau_s)}{A(\tau_s)} \right|^2 = 1$, $\left| \frac{B(-\tau_s)}{B(\tau_s)} \right|^2 = 0$.

FIG. 4. Probability distribution $W^{(n)}_{\text{ground}}(\nu)$ of definite-momentum states on ground level prepared asymmetrically by the interaction with the standing wave. Integer $n$ (as in Fig. 2 and 3) determines the value of momentum ($p = n\hbar k$). Momentum distribution before the travelling wave was symmetric. The chosen parameters are $2\nu\tau_s = 40$, consequently $\left| \frac{A(-\tau_s)}{A(\tau_s)} \right|^2 = 1/2$, $\left| \frac{B(-\tau_s)}{B(\tau_s)} \right|^2 = 1/2$, $\nu t = 20$. 

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FIG. 5. Probability distribution $W_{excited}^{(n)}$ of definite-momentum states in excited level prepared asymmetrically by the interaction with the standing and travelling waves. The chosen parameters are the same as in Fig. 4.

FIG. 6. Temporal behavior of momentum of translational motion for ground energy level. Time interval includes all parts of interaction: with preparing-standing and final-travelling waves. All parameters have the same values as in fig. 2.

FIG. 7. Temporal behavior of momentum of translational motion for excited energy level. All parameters have the same values as in fig. 2.

FIG. 8. Temporal behavior of mean momentum of translational motion for ground energy level. Time interval includes all parts of interaction: with preparing-standing and final-travelling waves. All parameters have the same values as in fig. 2.

FIG. 9. Temporal behavior of mean momentum of translational motion for excited energy level. Time interval includes all parts of interaction: with preparing-standing and final-travelling waves. All parameters have the same values as in fig. 2.

FIG. 10. The ground level, definite momentum state of an atom (zone 1) preliminary transforms into a superposition one by coherent interaction with a resonant standing wave (zone 2). The next interaction with the travelling wave leads to large-scale changes in atomic momentum distributions.
$W_{\text{ground}}^{(n)}$
$W^{(n)}_{\text{excited}}$

- ■ Standing
- ○ Accumulating
$W^{(n)}_{\text{excited}}$
$\langle n_e \rangle$
Standing
$\langle p_g \rangle$

Standing

Travel
$\langle p_e \rangle$

Standing

Traveling
Ground level definite-momentum atom

Standing wave

zone 1

zone 2

Resonant Kapitza-Dirac splitting

Travelling wave

zone 3

Stem-Gerlach type splitting