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ATOMIC VELOCITY SELECTION BY INTERFERENCE IN TWO-PHOTON IONIZATION

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This paper discusses a new scheme for generating quantum coherence between different degrees of freedom of an atom interacting with two modes of the electromagnetic field. The presence of quantum interference in a two-photon coupling between the ground state of the atom and the continuum through two quasi-resonant intermediate states induces selective ionization of the atoms for particular combinations of the different parameters characterizing the degrees of freedom of the system, leading to quantum coherence between the internal state, the center-of-mass motion of the atom, and the electromagnetic field. The application of this method to the selection of an atomic velocity class is discussed.

1. Introduction

In recent years, considerable progress has been made in the manipulation of microscopic systems by taking advantage of their quantum properties. Two of the most impressive realizations are the demonstrations of "sub-quantum noise" measurements of electromagnetic quantities using non-classical coherence properties of the light field (squeezing[1] and Twin-photon experiments[2]), and atomic manipulation and cooling. Atom vapors were cooled to extremely low temperatures of the order of a few nanoKelvins, allowing the first experimental observation of the Bose-Einstein (degenerate gas) condensation.[3] Mechanisms relying on the high sensitivity of the atomic state to motion in laser waves exhibiting gradients of polarization or alignment attained temperatures close to the "photon-recoil energy limit", typically of a few microKelvins.[4] This limit being related to the random nature of the spontaneous emission, it seemed for a little while to be the ultimate one. However, new schemes have been devised, and experimentally demonstrated, to overcome this limit through suppression of spontaneous emission by quantum interference (velocity selection by coherent population trapping - VSCPT[5]), by velocity selection through a Raman process,[6] or by selective evaporation out of a magnetic trap[7] (this latter method, not relying on fluorescence cycles, is not limited by the photon-recoil effect). This paper discusses, with emphasis on the basic physical processes, a recently proposed method for generating quantum coherence thanks to quantum interference in the two-photon ionization of atoms, that shows up to be able to generate non-classical light coherence[8] and to allow velocity selection of atoms.[9]

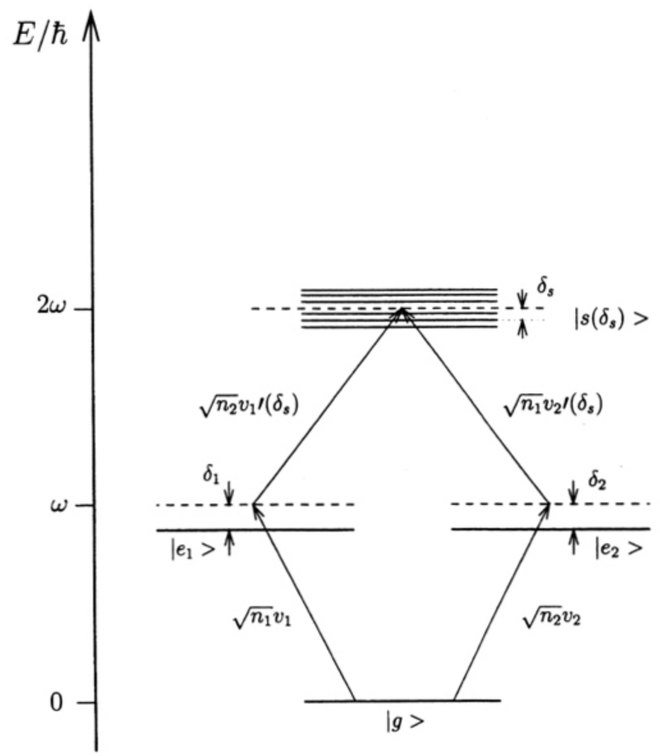

Fig. 1. Level scheme and electromagnetic couplings of the two-photon ionization scheme.

2. Interference Effects in Two-Photon Ionization

Consider a system where the ground state is coupled to the continuum by two distinct modes of the electromagnetic field through two quasi-resonant intermediate states (see Fig. 1). Mode 1, containing n_1 photons, couples the ground state g to the intermediate state e_1 (dipole matrix element v_1) and the intermediate state e_2 to the continuum (dipole matrix element v_1'). Mode 2 (n_2) couples the ground state g to the intermediate state e_2 (v_2) and the intermediate state e_2 to the continuum (v_2'). As there are two distinct paths leading from the ground state to a particular final state in the continuum, the two-photon ionization process shows interference effects. Let us recall that the two-photon ionization rate of the ground-state is given by (we restrict ourselves here to the one-dimensional case)

$$\Gamma_g(v;n_1,n_2) = \Gamma_0\omega^2 n_1 n_2 \left| \frac{1}{\delta_1 + k_1 v + b_{11}n_1 + b_{12}n_2} - \frac{(-v_2 v_2)/(v_1 v_1)}{\delta_2 + k_2 v + b_{21}n_1 + b_{22}n_2} \right|^2 . \tag{1}$$

In writing this equation, we took into account two perturbing effects affecting the energy interval between the ground-state and the intermediate levels, to which the ionization rate is highly sensitive. The two denominators in Eq. (1) include, aside of the laser atom detuning δ_i, the first-order Doppler effect correction $k_i v$ (k_i is the wave number for mode i and v the center-of-mass velocity) and the light-shift correction $b_{ij}n_j$ induced by the light intensity in mode j. The light-shift coefficients b_{ij} are calculated[8] by application of second order perturbation theory. Finally, Γ_0 is a constant depending on other parameters of the system. We neglect, in Eq. (1), eventual processes connecting the ground-state to the continuum by absorption of two photons of the *same* mode; they do not show quantum interference and are not interesting for our purposes here (for a detailed discussion, see Ref. 8).

We now discuss the generation of quantum coherence. We first made the simplifying assumption that the dipole matrix elements are such that $(-v_2 v_2'/v_1 v_1) = 1$ in Eq. (1). It is then clear that if we take $\delta_1 = \delta_2 = \delta$, the destructive interference between the two possible ionization paths will cause the ionization rate to vanish in the absence of the perturbing effects. In the presence of these effects, the ionization rate still vanishes if the following (destructive interference) condition is satisfied:

$$k_1 v + b_{11}n_1 + b_{12}n_2 = k_2 v + b_{21}n_1 + b_{22}n_2 . \tag{2}$$

Supposing that we prepare the system in an initial global state such that the internal atomic state is the ground state, the external atomic state is a superposition of eigenstates of the center-of-mass velocity operator $V = P/M$ and each field mode is a superposition of photon-number states:

$$|\Psi_0\rangle = |g\rangle \otimes \sum_v c_v|v\rangle \otimes \sum_n a_n|n\rangle \otimes \sum_m b_m|m\rangle = |g\rangle \otimes \sum_{v,n,m} a_n b_m c_v |v;n,m\rangle \tag{3}$$

where the three summation symbols apply respectively to the velocity eigenstates, the number states in mode 1, and the number states in mode 2. Note that there is no quantum coherence between the internal, the external atomic state and the field modes.

If we let the system evolve for a time t, the survival probability for each state in the right-hand side of Eq. (3) is given by:

$$P_g(v;n,m;t) = |a_n|^2|b_m|^2|c_v|^2 \exp[-\Gamma_g(v;n,m)t] . \tag{4}$$

If t is long compared to the typical coherence-building time $t_c = [\Gamma_0\omega^2\langle n\rangle\langle m\rangle/\delta^2]^{-1}$, this probability will be negligible except for those states corresponding to values of

v, n, m satisfying the destructive interference condition [Eq. (2), for which thus $\Gamma_g = 0$], that we shall call the "uncoupled" states (in the sense that they are not coupled to the continuum). The state of the system at a time $t \gg t_c$ is thus of the form

$$|\psi\rangle = |g\rangle \otimes |\varphi_g\rangle + |\Phi\rangle \tag{5}$$

where $|\varphi_g\rangle$ is a superposition containing only "uncoupled" states, whereas $|\Phi\rangle$ contains the states where the atomic internal state is a continuum state (we suppose that the detuning δ is large enough to allow us to neglect the population of the intermediate states), and thus $\langle g|\Phi\rangle = 0$. One sees that the state described by Eq. (5) is not factorable, and that the selective ionization process has generated a complicate kind of quantum coherence involving the ground-state — that is correlated only to "uncoupled" states — the external atomic state (center-of-mass motion) and the field modes.

As a first example, let us consider the case where one can neglect the Doppler effect compared to the light-shift (the atoms are supposed to be cold enough), Eq. (2) then implies that the ground state becomes correlated to the field modes satisfying:

$$n_2 = \alpha n_1 + \beta \tag{6}$$

where α and β are related to the light-shift coefficients. If a measurement of the atomic state is made, giving as a result the ground state, the wave packet is reduced to the first term in the right-hand side of Eq. (5), and the correlation is completely "transferred" to the field. The inconvenience of this scheme is that, as the correlation increases, the probability of obtaining g as a result of the atomic state measurement decreases, and this is the major limitation of the present method.

In the next two sections, we shall discuss the opposite limit of negligible light-shift compared to the Doppler effect, that allows atomic velocity class selection.

3. Atomic Velocity Selection by Selective Ionization

In this section we consider the limit where the light-shifts in Eq. (2) are negligible compared to the Doppler effect. We take the field modes to be counterpropagating: $k_1 = -k_2 = k$, implying that the "uncoupled" states now satisfy

$$\delta_1 + kv = \delta_2 - kv\,. \tag{7}$$

We thus see, after the reasoning of the previous section, that the ground state will be correlated to the velocity eigenstates characterized by $v = (\delta_2 - \delta_1)/2k$. Furthermore, if we choose $\delta_2 = \delta_1 = \delta$, the zero-velocity class will be selected.[a] There is no need, in the present case, for an explicit measurement of the atomic internal state: the ionized atoms, not interacting with the lasers, eventually escape

[a]Care should be taken in talking about "velocity classes", as we are considering quantum states that are eigenstates of the velocity operator. However, it can be shown that in the present case the concepts of velocity eigenstate and velocity class are equivalent.[8]

the interaction region, and this takes the role of an implicit atomic state measurement. We also note that the atomic level and field modes configuration used here is exactly that of an one-dimensional optical molasses or of a magneto-optical trap. The two systems can furthermore be advantageously combined, as we will see later. For that reason, we call our arrangement an "Ionization-assisted Optical Molasses" (IAOM).

If the only absorption process were the two-photon ionization, the cooling is not limited to the photon recoil energy, as the ionization is irreversible and thus not followed by spontaneous emission. As the number of ground state atoms inside the IAOM diminishes, the velocity dispersion becomes closer and closer to zero, leading to a vanishing temperature (!) but also to a vanishing number of atoms (!). For reasonable values of the detuning, however, one cannot neglect the possibility for an atom to absorb one photon, get to an intermediate state and come back to the ground state by spontaneous emission, performing a fluorescence cycle, exactly as in usual Optical Molasses. If the temperature is larger than the equilibrium temperature in the Molasses (the "Doppler limit temperature" T_D, of the order of several tenths of microKelvins) this process contributes to cooling the atoms,[b] but if the ionization process leads the system to temperatures *below* the Doppler limit, the fluorescence cycles tends on the contrary to *heat* the system. In other words, there is competition between the two processes.

In order to study this competition, we wrote a kinetic equation describing the evolution of the velocity distribution $n(v,t)$. As one could have expected, this equation is of the form of the usual Fokker-Planck equation describing the Doppler cooling process[10] plus an additional term describing the ionization[9]:

$$\frac{\partial n}{\partial \tau} = \frac{\partial (Vn)}{\partial V} + \frac{\partial^2 n}{\partial V^2} - gV^2 n\,. \tag{8}$$

In writing this equation, we introduced dimensionless, scaled variables: τ is the time scaled to the dumping time τ_v of the velocity by the Doppler process, $V = v/v_D$ is the velocity normalized to the Doppler limit velocity $v_D = \sqrt{k_B T_d/M}$ and gV^2 is the normalized ionization rate expanded up to the order v^2: $\Gamma_g = g(v/v_D)^2\tau_v^{-1} + o(v^4)$. This equation has an analytical solution. We discuss here only its most relevant features, a detailed discussion being available from Ref. 9.

Supposing that we start the ionization from the Doppler limit, Eq. (8) leads to a number of atoms that asymptotically tends to zero as[c] $\exp[-2g^{-1/2}\tau]$, but to a stationary temperature $T_{st} = g^{-1/2}T_D$. We thus see the velocity selection is able to cool the system below the Doppler limit temperature, still at the price of an important decreasing in the number of atoms.

In order to eliminate this last inconvenient, we have also considered the possibility of applying the ionization to a system where a source continuously brings into

[b] For negative detunings.

[c] For $g \gg 1$.

the system new ground-state, Doppler temperature, atoms. The source effect can be taken into account by adding to Eq. (8) a source term:

$$\frac{\partial n}{\partial \tau} = \frac{\partial (Vn)}{\partial V} + \frac{\partial^2 n}{\partial V^2} - gV^2 n + \frac{s}{\sqrt{2\pi}} e^{-v^2/2}, \tag{9}$$

where s is the scaled source intensity (the atom flux per unit dumping time τ_v). We have not found an analytical solution for this equation, but a combined analytical and numerical analysis[9] shows that it still leads to a sub-Doppler stationary temperature, proportional to $g^{-1/3}$, and to a stationary number of ground state atoms proportional to $(sg^{-2/3})$. It can also be shown[9] that the phase-space atomic density is asymptotically independent of the temperature, it depends only on the ratio of the source to the ionization rate, so that one can have high densities even for very small temperatures if the source term is high enough. It is however necessary to keep in mind that Eqs. (8) and (9) are obtained in the so-called "limit of small jumps", which states that the atom momentum change in an elementary process (here, a fluorescence cycle) should be negligible compared to the momentum of the atom. As a consequence, the approach presented here is not valid for temperatures comparable to, or smaller than, the photon-recoil temperature.

4. Generalization to Ground-State Degenerate Atoms

In this section, we propose a scheme that would allow one to combine the velocity selection method presented in the previous section with Sisyphus- or alignment-type mechanisms (and even with VSCPT).[4,11]

Sisyphus or alignment mechanisms are available for ground-state degenerated atoms. We consider a system whose levels have angular momenta $J_g = 1$, $J_e = 1$ and $J_s = 0$ (see Fig. 2). This system thus presents Sisyphus/alignment mechanisms, and the ionization process can start from a temperature a few times larger than the photon-recoil temperature. The reason we used in Fig. 2 a discrete upper level s which is itself coupled to the continuum is to avoid parasitic ionizations corresponding to the absorption of two photons of the same mode,[d] not subjected to interference and thus inducing an additional, useless, loss of atoms. It can be shown[12] that for a sufficiently high coupling between the upper level and the continuum, the dynamics of this system is equivalent to the previous one.

We suggest here a two-dimensional scheme where ground-state atoms in the Zeeman sublevel $m_g = 0$ interact with $\sigma^+ - \sigma^-$ counterpropagating waves along the x-axis, allowing cooling by the alignment mechanism along this axis. The atoms in the $m_g = \pm 1$ sublevels interact with orthogonally polarized (lin $\perp$ lin) counterpropagating waves along the y-axis and undertake Sisyphus effect in this dimension. Moreover, optical pumping pump atoms from one kind of sublevel to the other, allowing cooling in two dimensions. When the system arrives at the equilibrium temperature, the ionizing laser connecting the upper level to the continuum is

[d] If s is a continuum state, two-photon processes such as $(J_g = 1,\ m_g = -1) \to (J_e = 1,\ m_e = 0) \to (J_s = 1,\ m_s = 1)$ can not be avoided.

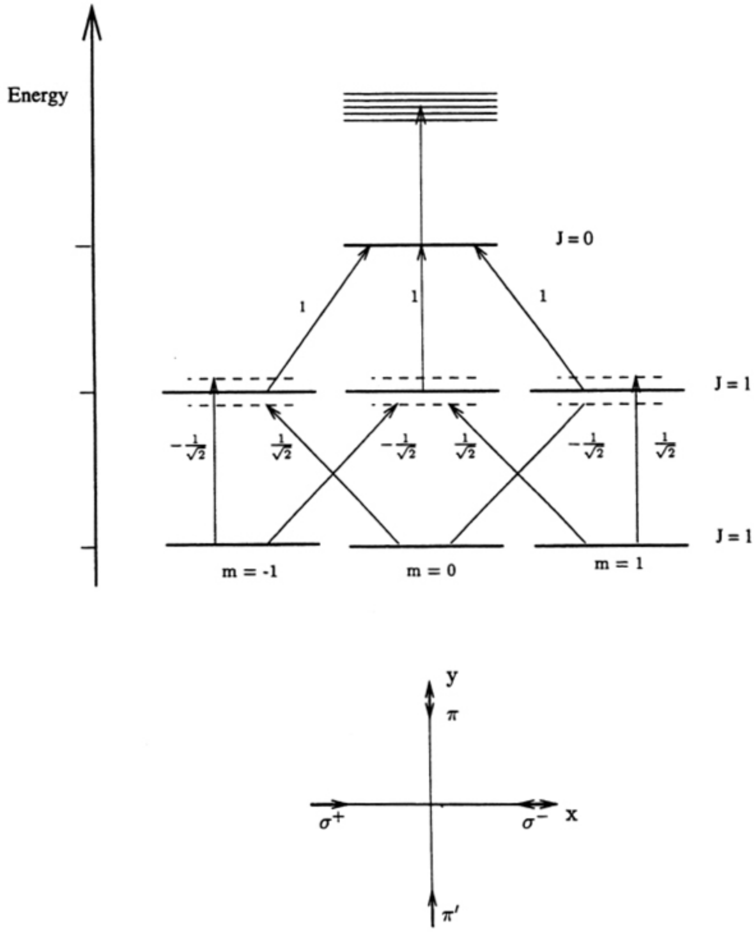

Fig. 2. $J = 1 \to J = 1 \to J = 0$ scheme for the two-dimensional selective ionization process allowing combination with Sisyphus-type mechanism.

turned on. For $m_g = 0$ atoms, the structure of the two-photon coupling is the same as in the previous section, it thus selects atoms with $v_x \approx 0$, for equal detunings $(\delta_{\sigma^+} = \delta_{\sigma^-})$. For $m_g = \pm 1$ atoms, a simple analysis shows that they are connected to the upper level by four different paths, leading to two possible final sates (see Fig. 3). The corresponding ionization rate is thus of the form (we take $m_g = 1$ for definiteness):

$$\Gamma_g(m_g = 1) = \overline{\Gamma}\omega^2 \left[\frac{n_{\sigma^-} - n_\pi}{2} \left| \frac{1}{\delta_\pi - k_\pi v_y} + \frac{1}{\delta_{\sigma^-} - k_{\sigma^-} v_x} \right|^2 + \frac{n_{\sigma^-} n_{\pi'}}{2} \left| \frac{1}{\delta_{\pi'} + k_{\pi'} v_y} + \frac{1}{\delta_{\sigma^-} - k_{\sigma^-} v_x} \right|^2 \right] \tag{10}$$

where the first two terms correspond to the absorption of a photon σ^- (propagating

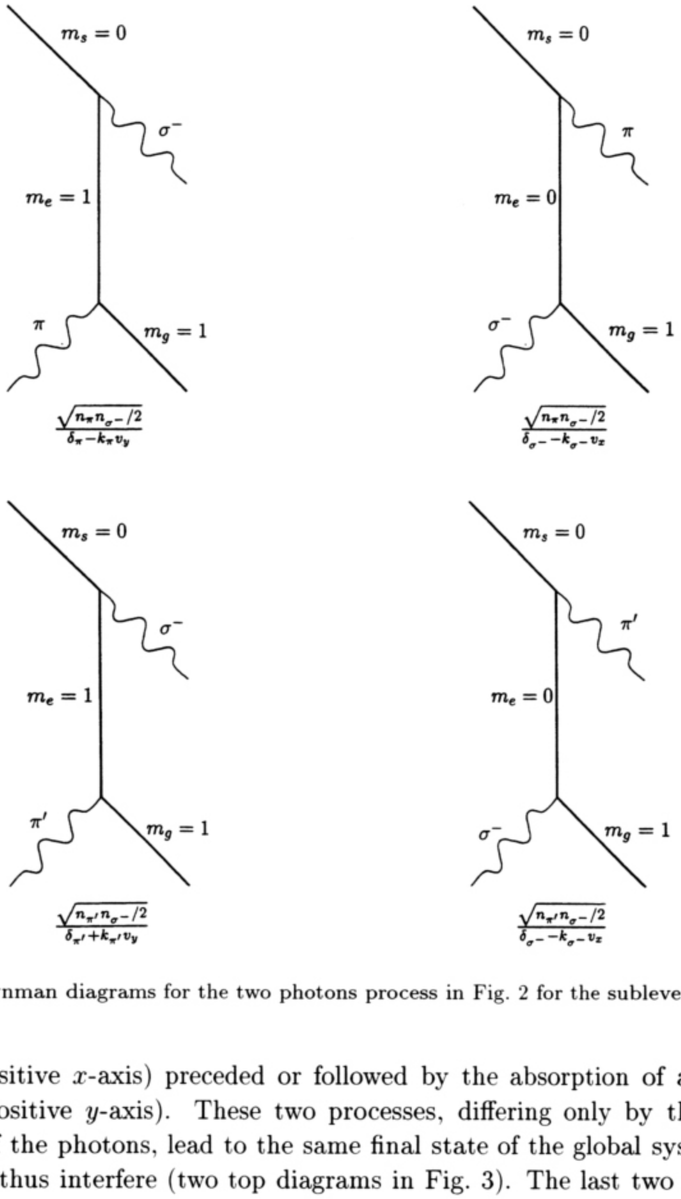

Fig. 3. Feynman diagrams for the two photons process in Fig. 2 for the sublevel $m_g = 1$.

along the positive x-axis) preceded or followed by the absorption of a photon π (along the positive y-axis). These two processes, differing only by the order of absorption of the photons, lead to the same final state of the global system (atom + field) and thus interfere (two top diagrams in Fig. 3). The last two terms (two bottom diagrams in Fig. 3) are of the same form, but the linearly polarized photon absorbed is a π' one, so that they do not lead to (and thus do not interfere with) the final state corresponding to the two first terms. Now, in order to experience destructive interference in the two-photon couplings, the low-velocity atoms must satisfy simultaneously *two* conditions, involving two degrees of freedom (v_x, v_y). A simple calculation shows that these conditions are compatible, provided $\delta_{\pi'} = \delta_\pi = -\delta_{\sigma^+}$. In this case, destructive interference is obtained for $v_x = v_y = 0$. An analogous calculation produces the same condition for $m_g = -1$ atoms. Again, optical pumping finishes by putting $m_g = 0$ atoms in one of the other sublevels, so that the two degrees of freedom of all atoms are affected by the two-photon process.

If the ionization process is supposed to begin from an equilibrium temperature times above the photon-recoil temperature,[11] roughly 10% of the atoms will be able to attain the photon-recoil limit, the rest being lost by ionization. This system seems thus a good candidate to produce sub-recoil temperatures. However, a detailed theory of the interaction between the selective ionization and the Sisyphus/alignment-type mechanism is still to be developed. We are presently working in a semiclassical version of this theory.[12]

Compared to VSCPT, that also relies on quantum interference, the scheme proposed here presents one major inconvenience: the decrease in the number of available atoms (except for the source version), whereas in VSCPT the number of trapped atoms in principle monotonically increases with the interaction time. On the other hand, an advantage of the IAOM is that the cooled atoms are in one of the sublevels of the ground state, and not in a coherent superposition. In consequence, this cooling process should be much less sensitive to collisions and other external perturbations.

5. Conclusion

In this paper, we discussed the possibility of applying selective ionization of atoms due to quantum interference to atomic velocity selection. We showed that the method can be advantageously combined with the usual Doppler cooling, and we suggested a scheme allowing to combine it with two-dimensional Sisyphus-alignment mechanisms, that should to be able to produce sub-recoil temperatures.

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