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Manipulating Atoms with Photons

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Abstract

By using quasi-resonant exchanges of energy, linear and angular momentum between atoms and photons, it is possible to polarize atoms, to displace their energy levels and to control their position and their velocity. A few physical mechanisms allowing one to trap atoms and to cool them in the microkelvin, and even in the nanokelvin range, are briefly described and classified with the help of a few simple guidelines. Various possible applications of such ultracold atoms are also reviewed. They take advantage of the long interaction times and long de Broglie wavelengths which are now available with laser cooling and trapping techniques. New quantum situations can also be achieved in these experiments, calling for new theoretical approaches. The last part of this paper is devoted to a brief discussion of these problems.

1. Introduction

Electromagnetic interactions play a central role in low energy physics. They are responsible for the cohesion of atoms and molecules and they are at the origin of the emission and absorption of light by such systems. There is another important application of atom-photon interactions, which will be discussed in this paper. Atom-photon interactions can be used to act on atoms, to manipulate them, to control their various degrees of freedom. In fact, this research field has considerably expanded during the last few years and it would be impossible to fully cover it in a short presentation. We will just try here to present a few guidelines for understanding these developments, for establishing connections between different physical effects. Since this symposium deals with basic quantum concepts and phenomena, we will also address a few important questions such as: What can one learn from these studies? What kind of new situations can be achieved? Can they improve our understanding of quantum phenomena?

The paper is organized as follows. We first present in Section 2 a brief review of basic physical mechanisms, trying to classify the various physical effects into two great categories: dissipative effects and dispersive effects. Such a partition applies to both internal and external degrees of freedom. Interesting effects can also result from the interplay between these two types of physical processes. We then briefly discuss in Section 3 important applications of ultracold atoms which result from the fact that they provide very long interaction times and very long de Broglie wavelengths. Finally, we mention in Section 4 a few problems, a few challenges which appear in connection with these developments. New quantum situations have been achieved which stimulate the development of new theoretical approaches. Great efforts have been also devoted to overcome the quantum limits, such as the single photon recoil limit in laser cooling.

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2. Brief review of physical mechanisms

2.1. *Existence of two types of effects in atom-photon interactions*

2.1.1. *Absorption and dispersion of light.* Consider a light beam with frequency ω_L propagating through a medium consisting of atoms with resonance frequency ω_A. The index of refraction describing this propagation has an imaginary part and a real part which are associated with two types of physical processes. The incident photons can be absorbed, more precisely scattered in all directions. The corresponding attenuation of the light beam is maximum at resonance. It is described by the imaginary part of the index of refraction which varies with $\omega_L - \omega_A$ as a Lorentz absorption curve. We will call such an effect a dissipative effect. The speed of propagation of light is also modified. The corresponding dispersion is described by the real part n of the index of refraction whose deviations to 1, $n - 1$, vary with $\omega_L - \omega_A$ as a Lorentz dispersion curve. We will call such an effect a dispersive effect.

2.1.2. *Broadening and shift of atomic energy levels.* Dissipative effects and dispersive effects also appear for the atoms, as a result of their interaction with photons. They correspond to a broadening and to a shift of the atomic energy levels, respectively.

These effects already appear when the atom interacts with the quantized radiation field in the vacuum state. It is well known that atomic excited states get a natural width Γ, which is also the rate at which a photon is spontaneously emitted from such states. Atomic energy levels are also shifted as a result of virtual emissions and reabsorptions of photons by the atom. Such a radiative correction is nothing but the Lamb shift [1].

Similar effects are associated with the interaction with an incident light beam. Atomic ground states get a radiative broadening Γ', which is also the rate at which photons are absorbed, more precisely scattered, from the incident beam. Atomic energy levels are also shifted as a result of virtual absorptions and reemissions of the incident photons by the atom. Such displacements $\hbar\Delta'$ are called light shifts, or ac Stark shifts [2, 3].

In view of their importance for the following discussions, we give now a brief derivation of the expressions of Γ' and Δ', using the so-called dressed-atom approach (see for example [4], chapter 6). In the absence of coupling ($V_{AL} = 0$, where V_{AL} is the atom-light interaction Hamiltonian), the energy levels of the "atom + photons" system bunch into well separated two dimensional manifolds $\{|g, N\rangle, |e, N-1\rangle\}$ where $|g, N\rangle$ (resp. $|e, N-1\rangle$) represents an atom in the ground state g (resp. excited state e) in the presence of N (resp. $N - 1$) photons. The left part of the Fig. 1 represents two such states, which are separated by an energy

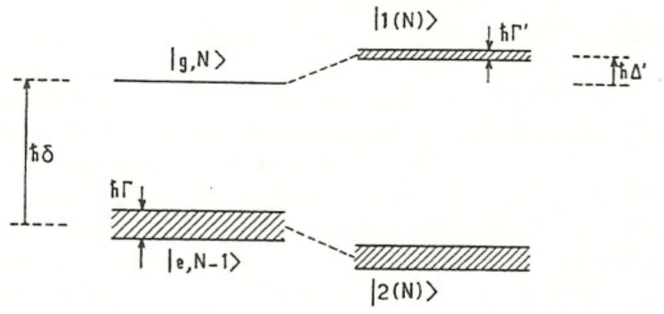

Fig. 1. Left part – Non coupled states of the atom-photon system. $\delta = \omega_L - \omega_A$ is the detuning between the laser and atomic frequencies, Γ the natural width of the excited state e. Right part – Dressed states appearing as a result of atom-photon interactions. Γ' is the radiative broadening of the ground state g, $\hbar\Delta'$ the light shift.

interval $\hbar\delta$ where $\delta = \omega_L - \omega_A$ is the detuning between the light frequency ω_L and the atomic frequency ω_A. We have represented by hatched lines the energy width $\hbar\Gamma$ of the excited state. Mathematically, this is described by an imaginary part, $-i\hbar\Gamma/2$, added to the energy of e. The interaction Hamiltonian V_{AL} couples $|g, N\rangle$ to $|e, N-1\rangle$, because the atom in g can absorb one photon and jump to e. The corresponding matrix element is equal to $\hbar\Omega_1/2$, where Ω_1 is the so-called Rabi frequency proportional to $\sqrt{N}$ and to the atomic dipole moment. Such a coupling gives rise to the two dressed states represented in the right part of Fig. 1 and which are obtained by diagonalizing the effective non-Hermitian Hamiltonian:

$$H_{\text{eff}} = \hbar\begin{pmatrix} \delta & \Omega_1/2 \\ \Omega_1/2 & -i\Gamma/2 \end{pmatrix}. \quad (1)$$

The energy $E_{g,N}$ of $|g, N\rangle$ is transformed into:

$$\tilde{E}_{g,N} = E_{g,N} + \hbar\Delta' - i\hbar\frac{\Gamma'}{2}. \quad (2)$$

The contamination of $|g, N\rangle$ by $|e, N-1\rangle$, due to V_{AL}, transfers to $|g, N\rangle$ part of the radiative instability of $|e, N-1\rangle$. This is the origin of the radiative broadening Γ'. The repulsion between the two levels explains the light shift Δ'. Simple expressions can be derived for Γ' and Δ' in the weak intensity limit where Ω_1 is small compared to Γ or $|\delta|$. A perturbative calculation of the eigenvalues of (1) gives:

$$\Gamma' = \Omega_1^2\frac{\Gamma}{\Gamma^2 + 4\delta^2}, \quad (3)$$

$$\Delta' = \Omega_1^2\frac{\delta}{\Gamma^2 + 4\delta^2}. \quad (4)$$

Both Γ' and Δ' are proportional to $\Omega_1^2 \propto N$, i.e. to the light intensity. They vary with the detuning $\delta = \omega_L - \omega_A$ as Lorentz absorbtion and dispersion curves, respectively, which justifies the denominations dissipative and dispersive used for these two types of effects. For large detunings ($|\delta| \gg \Gamma$), Γ' varies as $1/\delta^2$ and becomes negligible compared to Δ' which varies as $1/\delta$. On the contrary, for small detunings ($|\delta| \ll \Gamma$), Γ' is much larger than Δ'. In the high intensity limit, when Ω_1 is large compared to Γ and $|\delta|$, the two dressed states are the symmetric and antisymmetric linear combinations of $|g, N\rangle$ and $|e, N-1\rangle$. Their splitting is $\hbar\Omega_1$ and they share the instability Γ of e in equal parts, so that $\Gamma' = \Gamma/2$. One can explain in this way various physical effects such as the Rabi flopping or the Autler-Townes splittings of the spectral lines connecting e or g to a third level [5].

2.2. *Manipulation of internal degrees of freedom*

2.2.1 *Optical pumping.* Optical pumping is one of the first examples of manipulation of atoms by light [6]. It uses resonant excitation of atoms by circularly polarized light for transferring to the atoms part of the angular momentum carried by the light beam. It is based on the fact that different Zeeman sublevels in the atomic ground state have in general different absorption rates for an incoming polarized light. For example, for a $J_g = 1/2 \leftrightarrow J_e = 1/2$ transition, only atoms in the sublevel $M_g = -1/2$ can absorb σ^+-polarized light. They are excited in the sublevel $M_e = +1/2$ of e from which they can fall back in the sublevel $M_g = +1/2$ by spontaneous emission of a π-polarized photon. They remain then trapped in this state because no σ^+ transition starts from it. It is possible in this way to obtain high degrees of polarization in atomic ground states.

2.2.2. *Light shifts.* Optical pumping is a dissipative effect because it is associated with resonant absorption of photons by the atom. Non resonant optical excitation produces light shifts of the ground state Zeeman sublevels. Because of the polarization selection rules, light shifts depend on the polarization of the exciting light and vary in general from one Zeeman sublevel to another. For example, for a $J_g = 1/2 \leftrightarrow J_e = 1/2$ transition, a σ^+-polarized excitation shifts only the Zeeman sublevel $M_g = -1/2$, whereas a σ^--polarized excitation shifts only the sublevel $M_g = +1/2$. Magnetic resonances curves in the ground state g, which are very narrow because the relaxation time in g can be very long, are thus light shifted by a polarized non resonant excitation. It is in this way that light shifts were first observed [3, 7].

Light shifts can be considered from different points of view. First, they can be interpreted as a radiative correction, due to the interaction of the atom with an incident field rather than with the vacuum field. This is why Alfred Kastler was calling them "Lamp shifts". Secondly, they introduce perturbations to high precision measurements using optical methods, which must be taken into account before extracting spectroscopic data from these measurements. Finally, they are now more and more frequently used for manipulating the energy of Zeeman sublevels and for producing spatial modulations of these energies on an optical wavelength scale, which would not be easily achieved with spatially varying magnetic fields. We will see in Section 2.3.4 interesting applications of such a situation.

2.3. *Manipulation of external degrees of freedom*

2.3.1. *The two types of radiative forces.* There are two types of radiative forces associated with dissipative effects and reactive effects, respectively (for a review of radiative forces, see [8]).

Dissipative forces, also called radiation pressure forces or scattering forces, are associated with the transfer of linear momentum from the incident light beam to the atom in resonant scattering processes. They are proportional to the scattering rate Γ'. Consider for example an atom in a laser plane wave with wave vector $\boldsymbol{k}$. Because photons are scattered with equal probabilities in two opposite directions, the mean momentum transferred to the atom in an absorption-spontaneous emission cycle is equal to $\hbar k$. The mean rate of momentum transfer, i.e. the mean force, is thus equal to $\hbar k\Gamma'$. Since Γ' saturates to $\Gamma/2$ at high intensity (see Section 2.1.2), the radiation pressure force saturates to $\hbar k\Gamma/2$. The corresponding acceleration (or deceleration) a which can be communicated to an atom with mass M, is equal to $a_{\max} = \hbar k\Gamma/2M = v_R/2\tau$, where $v_R = \hbar k/M$ is the recoil velocity of the atom absorbing or emitting a single photon, and $\tau = 1/\Gamma$ is the radiative lifetime of the excited state. With $v_R \simeq 10^{-2}$ m/s and $\tau \simeq 5.10^{-9}$ s, a can reach values as large as 10^6 m/s^2, i.e. 10^5 g where g is the gravity acceleration. With such a force, one can stop a thermal atomic beam in a distance of the order of one meter, provided that one compensates for the Doppler shift of the decelerating atom, by using for example a spatially varying Zeeman shift [9, 10].

Dispersive forces, also called dipole forces or gradient forces, are associated with position dependent light shifts $\hbar\Delta'(\boldsymbol{r})$ due to a spatial varying light intensity. Consider for example a laser beam largely detuned from resonance, so that one can neglect Γ' (no scattering process). The atom thus remains in the ground state and the light shift $\hbar\Delta'(\boldsymbol{r})$ of this state plays the role of a potential energy, giving rise to a force which is equal to the opposite of its gradient: $\boldsymbol{F} = -\nabla[\hbar\Delta'(\boldsymbol{r})]$. Such a force can be also interpreted as resulting from a redistribution of photons between the various plane waves forming the laser wave in absorption-stimulated emission cycles. If the detuning is not large enough to allow Γ' to be neglected, spontaneous transitions occur between dressed states having opposite gradients, so that the instantaneous force oscillates back and forth between two opposite values in a random way. Such a dressed atom picture provides a simple interpretation of the mean value and of the fluctuations of dipole forces [11].

2.3.2. *Applications of dissipative forces – Doppler cooling and magneto-optical traps.* We have already mentioned in the previous subsection the possibility to decelerate an atomic beam by the radiation pressure force of a laser plane wave. Interesting effects can also be obtained by combining the effects of two counterpropagating laser waves.

A first example is Doppler cooling [12, 13] which results from a Doppler induced imbalance between two opposite radiation pressure forces. The two counterpropagating laser waves have the same (weak) intensity and the same frequency and they are slightly detuned to the red of the atomic frequency ($\omega_L < \omega_A$). For an atom at rest, the two radiation pressure forces exactly balance each other and the net force is equal to zero. For a moving atom, the apparent frequencies of the two laser waves are Doppler shifted. The counterpropagating wave gets closer to resonance and exerts a stronger radiation pressure force than the copropagating wave which gets farther from resonance. The net force is thus opposite to the atomic velocity v and can be written for small v as $F = -\alpha v$ where α is a friction coefficient. By using three pairs of counterpropagating laser waves along three orthogonal directions, one can damp the atomic velocity in a very short time, on the order of a few microseconds, achieving what is called an "optical molasses" [14]. The theory of Doppler cooling predicts that the equilibrium temperature obtained with such a scheme is always larger than a certain limit T_D, called the Doppler limit, and given by $k_B T_D \sim \hbar\Gamma$ where Γ is the natural width of the excited state and k_B the Boltzmann constant. Such a limit, which is reached for $\delta = \omega_L - \omega_A = -\Gamma/2$, is, for alkali atoms, on the order of 100 μK. In fact, when the measurements became precise enough, it appeared that the temperature in optical molasses was much lower than expected [15]. This indicates that other laser cooling mechanisms, more powerful than Doppler cooling, are operating. We will come back to this point in Section 2.3.4.

The imbalance between two opposite radiation pressure forces can be also made position dependent though a spatially dependent Zeeman shift produced by a magnetic field gradient. In a one-dimensional configuration, first suggested by J. Dalibard in 1986, the two counterpropagating waves, which are detuned to the red ($\omega_L < \omega_A$) and which have opposite circular polarizations are in resonance with the atom at different places. This results in a restoring force towards the point where the magnetic field vanishes. Furthermore the non zero value of the detuning provides a Doppler cooling. In fact, such a scheme can be extended to three dimensions and leads to robust, large and deep traps [16]. It combines trapping and cooling, it has a large velocity capture range and it can be used for trapping atoms in a cell [17].

2.3.3. *Applications of dispersive forces: laser traps and atomic mirrors.* When the detuning is negative ($\omega_L - \omega_A < 0$), light shifts are negative. If the laser beam is focussed, the focal zone where the intensity is maximum appears as a minimum of potential energy, forming a potential well where sufficiently cold atoms can be trapped. This is a laser trap. Laser traps using a single focussed laser beam [18] or two crossed focussed laser beams [19, 20] have been achieved.

If the detuning is positive, light shifts are positive and can thus be used to produce potential barriers. For example an evanescent blue detuned wave at the surface of a piece of glass can prevent slow atoms impinging on the glass surface from touching the wall, making them bouncing off a "carpet of light" [21]. This is the principle of mirrors for atoms. Plane atomic mirrors [22, 23] have been achieved as well as concave mirrors [24]. The advantage of concave mirrors is that the transverse atomic motion is stable if atoms are released from a point located below the focus of the mirror. It is then possible to observe several successive bounces of the atoms and such a system can be considered as a "trampoline for atoms" [24]. See Fig. 2.

2.3.4. *Combination of dispersive effects and dissipative effects: Sisyphus cooling.* When the detuning is not too large, both light shifts and optical pumping rates can have appreciable values and their combination can give rise to efficient laser cooling mechanisms that we describe now (for a review, see [25]). In the usual laser cooling configurations, the polarization of the resulting total laser field is in general spatially modulated. For example, it may change from σ^+ to σ^- (and vice versa) every $\lambda/4$, being elliptical or linear in between. Two ground state Zeeman sublevels $M_g = \pm 1/2$ (denoted $\pm$ in Fig. 3) have then in general different light shifts which are spatially modulated with the same period as the laser polarization. The spatial modulation of the laser polarization also gives rise to a spatial modulation of the optical pumping rates, atoms being pumped from $M_g = -1/2$ to $M_g = +1/2$ in the places where the polarization is σ^+, and from $M_g = +1/2$ to $M_g = -1/2$ in the places where

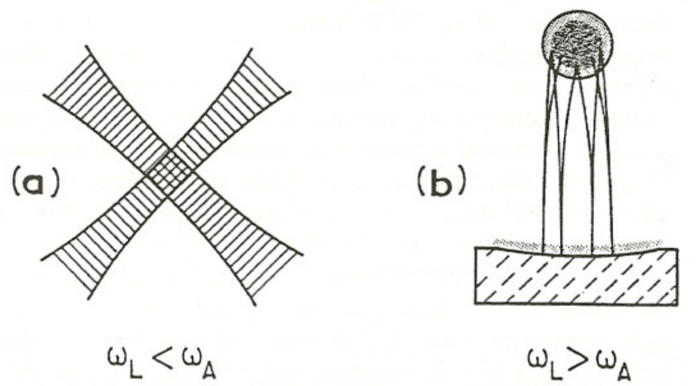

Fig. 2. (a) Laser trap appearing at the intersection of the focal zones of two laser beams with a negative detuning ($\delta = \omega_L - \omega_A < 0$). (b) Atomic mirror associated with an evanescent wave covering a glass surface. The detuning δ is now positive. If the glass surface is concave, the transverse motion of the atom can be stable and the system can be considered as a trampoline for atoms.

the polarization is σ^-. The two spatial modulations of light shifts and optical pumping rates are correlated because they have the same cause, the spatial modulation of the laser polarization. For a proper choice of the detuning, these correlations result in the fact that optical pumping always transfers atoms from the higher Zeeman sublevel to the lower one (see Fig. 3). This can lead to a very efficient cooling mechanism, called "Sisyphus cooling" or "polarization gradient cooling" [26, 27] (see also [11]). Consider an atom moving to the right and starting from the bottom of a valley, for example in the state $M_g = +1/2$ at a place where the polarization is σ^+. The atom can climb up the potential hill and reach the top of the hill where is has the maximum probability to be optically pumped in the other sublevel, i.e. in the bottom of a valley, and so on (see Fig. 3). Like Sisyphus in the Greek mythology, the atom is running up potential hills more frequently than down. When it climbs a potential hill, its kinetic energy is transformed into potential energy which is then dissipated by light, since the spontaneously emitted photon has an energy higher than the absorbed laser photon (anti-Stokes Raman processes of Fig. 3). Such a cooling mechanism is very efficient and can lead to temperatures T on the order of a few microkelvins, given by $k_B T \simeq U_0$, where U_0 is the depth of the optical potential wells of Fig. 3, i.e. the maximum differential light shift. Equation $k_B T \simeq U_0$ cannot remain valid when U_0 tends to zero, because we have neglected the recoil due to the spontaneously emitted photons. There is a threshold for U_0, on

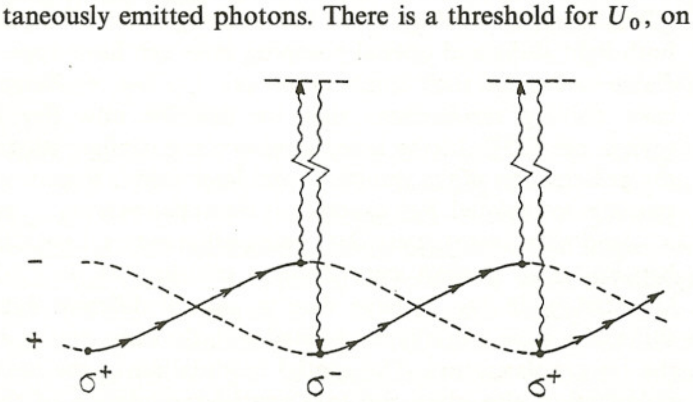

Fig. 3. Principle of Sisyphus cooling. Because of the correlations which exist between the spatial modulations of light shifts and optical pumping rates, the moving atom is running up potential hills more frequently than down.

the order of a few recoil energies $E_R = \hbar^2 k^2/2M$, below which Sisyphus cooling can no longer work.

Note finally that, for the optimal conditions of Sisyphus cooling, atoms become so cold that they get trapped in the quantum vibrational levels of Fig. 3. More precisely, one must consider energy bands in this perodic structure [28]. Experimental observation of such a quantization of atomic motion in an optical potential has been first achieved at one dimension [29, 30]. Atoms then form a spatial periodic array, called "1D-optical lattice", with an antiferromagnetic order, since two adjacent potential wells correspond to opposite spin polarizations. 2D and 3D optical lattices have been achieved subsequently (see the review papers [31, 32]).

3. Brief survey of applications

3.1. *Long observation times*

3.1.1. *Atomic clocks using ultracold atoms.* Ultracold atoms move very slowly. For example, Cesium atoms cooled by Sisyphus cooling have an effective temperature on the order of 1 μK, corresponding to a r.m.s. velocity of 1 cm/s. This allow them to spend a longer time T in an observation zone where a microwave field induces resonant transitions between the two hyperfine levels g_1 and g_2 of the ground state. Increasing T decreases the width $\Delta\nu \sim 1/T$ of the microwave resonance line whose frequency is used to define the unit of time. The stability of atomic clocks can thus be considerably improved by using ultracold atoms [33, 34].

In usual atomic clocks, atoms from a thermal cesium beam cross two microwave cavities fed by the same oscillator. The average velocity of the atoms is several hundred m/s, the distance between the two cavities is on the order of 1 m. The microwave resonance between g_1 and g_2 is monitored and is used to lock the frequency of the oscillator to the center of the atomic line. The narrower the resonance line, the more stable is the atomic clock. In fact, the microwave resonance line exhibits Ramsey interference fringes whose width $\Delta\nu$ is determined by the time of flight T of the atoms from one cavity to another. For the longest devices, T, which can be considered as the observation time, can reach 10 ms, leading to values of $\Delta\nu \sim 1/T$ on the order of 100 Hz.

Much narrower Ramsey fringes, with sub-Hertz linewidths can be obtained in the so-called "Zacharias atomic fountains". Atoms are captured in a magneto-optical trap and laser cooled before being launched upwards by a laser pulse through a microwave cavity. Because of gravity they are decelerated, they return and fall back, passing a second time through the cavity. Atoms therefore experience two coherent microwave pulses, when they pass through the cavity, the first time on their way up, the second time on their way down. The time interval between the two pulses can now be on the order of 1 s, i.e. about two order of magnitudes longer than with usual clocks. Atomic fountains have been achieved for sodium [35] and cesium [36]. A short-term relative frequency stability of $2 \times 10^{-13}\tau^{-1/2}$, where τ is the integration time, has been recently measured for a one meter high Cesium fountain [37]. For $\tau = 10^4$ s, this gives $\Delta\nu/\nu \sim 2 \times 10^{-15}$. In fact such a stability is probably limited by the Hydrogen maser which is used as a reference source and the real stability, which could be more precisely determined by beating the signals of two fountain clocks, is expected to be $\Delta\nu/\nu \sim 10^{-16}$ for a one day integration time.

To increase the observation time beyond one second, a possible solution consists of building a clock for operation in a reduced gravity environment. Tests have been performed with an experimental set up embarked in a plane making parabolic free flights. A compact ultra stable prototype has been also built which could work in a satellite and a 7 Hz resonance signal has been recorded in a $10^{-2}g$ environment.

3.1.2. *Fundamental tests.* Atomic clocks working with ultracold atoms could not only provide an improvement of the Global Positioning System (GPS). They could be also used for basic studies.

A first possibility could be to build two fountains clocks, one with Cesium and one with Rubidium, in order to measure with a high accuracy the ratio between the hyperfine frequencies of these two atoms. Because of relativistic corrections, the hyperfine splitting is a function of $Z\alpha$ where α is the fine structure constant and Z is the atomic number [38]. Since Z is not the same for Cesium and Rubidium, the ratio of the two hyperfine structures depends on α. By making several measurements of this ratio over long periods of time, one could check Dirac's suggestion concerning a possible variation of α with time. The present upper limit for $\dot{\alpha}/\alpha$ in laboratory tests could be improved by two orders of magnitude.

Another interesting test would be to measure with a higher accuracy the gravitational red shift and the gravitational delay of an electromagnetic wave passing near a large mass (Shapiro effect [39]).

3.2. *Long de Broglie wavelengths*

When the velocity v decreases, the de Broglie wavelength $\lambda_{dB} = h/mv$ increases and several effects related to the wave nature of atomic motion become important.

3.2.1. *Atom optics – Atom interferometry.* New research fields, such as atom optics and atom interferometry, have appeared during the last few years, extending to atomic de Broglie waves the various experiments which were previously achieved with electromagnetic waves (for a general review, see [40] and [32]).

In atom optics, atomic de Broglie waves are manipulated with the equivalent of optical elements such as lenses, mirrors, beam splitters . . . An example of mirror for atoms has been described above (Section 2.3.3). The possibility to introduce friction forces which damp the atomic velocity adds an extra degree of freedom to atom optics which can thus become dissipative: the etendue of an atomic beam is no longer conserved and "funnels" for atoms can be achieved. Important applications have been recently developed, such as atom lithography, which allows atoms to be deposited on a substrate to form controlled structures with a resolution of a few tens of nanometers.

Interferometers for atomic de Broglie waves have been demonstrated. In contrast to light interferometers, they are sensitive to gravitational effects which can be measured with a great accuracy. The equivalent of the Sagnac effect for light has been observed with atomic de Broglie waves (the most recent experiment is described in [41]). The inherent sensitivity of such "atomic gyrometers" can exceed that of photon gyrometers by a factor of $Mc^2/\hbar\omega \sim 10^{11}$ (Mc^2 is the rest mass energy of the atom, $\hbar\omega$ is the energy of the photon). Such a large factor is due to the fact that slow atoms spend a much longer time in the interferometer than photons.

3.2.2 *Two recent experiments demonstrating the wave nature of atomic motion.* In the trampoline for atoms described in Section 2.3.2, atoms are considered as classical particles bouncing off a concave mirror. The associated de Broglie waves are also reflected by the mirror and one can consider for such a "gravitational cavity" standing de Broglie waves analogous to the light standing waves for a Fabry–Perot cavity [42]. By modulating at frequency $\Omega/2\pi$ the intensity of the evanescent wave which forms the atomic mirror, one can produce the equivalent of a vibrating mirror for de Broglie waves. The reflected waves thus have a modulated Doppler shift. The corresponding frequency modulation of these waves has been recently demonstrated [43] by measuring the energy change ΔE of the bouncing atom, which is found to be equal to $n\hbar\Omega$, where $n = 0, \pm 1, \pm 2, \ldots$ The discrete nature of this energy spectrum is a pure quantum effect. For classical particles bouncing off a vibrating mirror, ΔE would vary continuously in a certain range.

As a second example of quantum effect we will mention the recent observation of Bloch oscillations of ultracold atoms in a periodic optical potential [44]. Such a potential is in fact the spatially modulated light shift produced by a non resonant light standing wave. By chirping the frequency of the two counterpropagating laser waves forming the standing wave, one can produce an accelerated standing wave. In the rest frame of this wave atoms thus feel a constant inertial force in addition to the periodic potential. They are accelerated and the de Broglie wavelength λ_{dB} decreases. When $\lambda_{dB} = \lambda_{Laser}$, the de Broglie wave is Bragg reflected by the periodic optical potential. Instead of increasing linearly with time, the mean velocity $\langle v \rangle$ of the atoms oscillates back and forth. Such Bloch oscillations, which are a well known effect of solid state physics, are more easily observed with ultracold atoms than with electrons in condensed matter because the Bloch period can be much longer than the time between collisions which destroy the coherence of de Broglie waves.

3.2.3. *Quantum degenerate gases.* At very low temperatures and high densities, the average distance between atoms can become comparable to the thermal de Broglie wavelength $\lambda_{dB} = \sqrt{2\pi\hbar^2/Mk_B T}$. Equivalently, the phase-space density $n\lambda_{dB}^3$, where n is the number of atoms per unit volume, can become larger than 1. In this regime, spectacular effects can occur as a consequence of the symmetry properties of the wave function describing an ensemble of identical particles. Bosons trapped in an external potential are predicted to condense in the quantum ground state of the trap. Up to now, the only evidence for such an effect, called Bose–Einstein Condensation (BEC), came from studies on superfluid liquid Helium and excitons in semiconductors. The strong interactions which exist in such systems modify qualitatively the nature of the condensation. A great challenge was therefore to observe BEC in an ultracold dilute atomic gas where interactions are much weaker.

During the last two decades, great efforts have been devoted to observe BEC in spin-polarized hydrogen, which is the only quantum gas to remain gaseous at absolute zero. New techniques for cooling and trapping atomic gases have been developed, such as evaporative cooling and magneto-static trapping (for recent reviews, see [45, 46, 47]). Phase space densities very close to 1 have been achieved.

It is only very recently that it has been possible, by combining laser manipulation techniques with evaporative cooling and magneto-static trapping, to observe BEC on alkali atoms, such as rubidium [48], sodium [49], lithium [50]. A new exciting research field is being opened by these experiments. For a recent review, we refer the reader to [51].

4. New quantum problems – New challenges

The improvement in our ability to manipulate atomic systems has led to new situations where quantum effects play an essential role in atom-photon interactions. The analysis of these situations has stimulated the development of new theoretical approaches and has, for example, provided a better physical insight into the quantum evolution of a single atomic system. A great challenge has been also to overcome the quantum limits, such as the single photon recoil limit in laser cooling. In this last section, we briefly discuss a few of these problems. We focus here on the atomic evolution since other papers of this symposium deal with the generation of non classical light.

4.1. *A few examples of quantum situations*

At very low temperatures, a semiclassical description of atomic motion in laser light is no longer valid. The coherence length of the atomic wave packets can no longer be neglected in comparison with the laser wavelength. A particularly interesting situation occurs when the atomic momentum spread δp is smaller than the photon momentum $\hbar k$ (subrecoil cooling – see Section 4.3). Condition $\delta p < \hbar k$ is equivalent to $\lambda_{\mathrm{dB}} = h/\delta p > \lambda_{\mathrm{laser}} = h/\hbar k$. The atomic de Broglie wave can then remain coherent over several laser wavelengths. A full quantum treatment of both internal and external atomic degrees of freedom becomes essential. In particular, the concept of radiative force in a given point has to be generalized because it is not compatible with Heisenberg's uncertainty relations for the atomic center of mass (for a general review of atomic motion in laser light, see [8]).

Most laser cooling schemes introduce a damping of the atomic momentum. This is a dissipative process, necessarily associated with fluctuations. In the problems studied here, these fluctuations are due to spontaneous emission of fluorescence photons which can be emitted at random times and in random directions. The interest of spontaneous emission is that it provides a simple basic model of a quantum dissipative process. Understanding and describing the effect of these fluctuations is a central problem of quantum optics. It has given rise to a lot of theoretical works using different approaches such as master equations, Heisenberg-Langevin equations, quantum Fokker–Planck equations, dressed-atom approach … (for a general review, see [52] or [4]). The fluctuations of the atomic evolution appear in a very spectacular way in single-atom experiments (for a general review, see [53]). For example, the fluorescence light emitted by a single trapped ion excited by a resonant laser beam can appear and disappear in a random way, indicating that the atom is performing "quantum jumps" between different energy levels [54–56]. The interpretation of these experiments has stimulated the development of new descriptions of quantum dissipative processes which we briefly describe now.

4.2. *A new physical picture: random walks in Hilbert space*

A single atom excited by a resonant laser beam emits a random sequence of fluorescence photons which can be interpreted in terms of a radiative cascade of the atom dressed by the laser photons. Such a point of view allows a simple calculation of the "delay function" $W(\tau)$ (also called the "waiting time distribution") giving the distribution of the time intervals τ between two successive spontaneous emissions [57]. Between two such jumps of the radiative cascade, the dressed atom evolves in a coherent way in a given manifold and this coherent evolution is governed by an effective Hamiltonian analogous to the one given by eq. (1). One can thus picture the atomic evolution as consisting of a sequence of quantum jumps occurring at random times and separated by coherent evolution periods. This leads to a simple interpretation of the experimental observations. In particular, one can use $W(\tau)$ to make Monte-Carlo simulations of the random sequence of quantum jumps [58].

More general situations, where the delay function cannot be easily calculated have been also investigated, using quantum trajectories [59] or Monte-Carlo wave functions ([60] and references therein). All these new descriptions of dissipative quantum optics problems lead to a new physical picture, consisting of a "stochastic wave function" performing a random walk in Hilbert space. One can show that such approaches lead to the same results as the usual master equation treatments using density matrices. They have however a certain number of advantages. First, they provide a possible description of the quantum dissipative evolution of a single atomic system. Second, they lead in general to more efficient numerical calculations because it is easier to deal with state vectors than with density matrices. Finally, they allow fruitful connections to be established with other areas of statistical physics such as Brownian motion or anomalous random walks.

4.3. *Example of quantum limit*

4.3.1. *The single photon recoil limit. How to circumvent it.* The random recoil communicated to the atom by a spontaneously emitted photon seems at first sight to introduce a fundamental limit in laser cooling. If fluorescence cycles never stop, one cannot reduce the atomic momentum spread δp below the value $\hbar k$ corresponding to the momentum of a single photon.

It is however possible to circumvent such a limit and to achieve what is called a "subrecoil" laser cooling. The basic idea is to create a situation where the photon absorption rate R (which is also the rate at which photons are then spontaneously reemitted) depends on the atomic velocity $v = p/M$ and vanishes for $v = 0$ [Fig. 4(a)]. Atoms with $v = 0$ no longer absorb light. They no longer recoil in a random way and they remain cold. The other atoms undergo fluorescence cycles which produce random changes of v. The corresponding random walk in v-space can bring them in a small zone around $v = 0$ where they remain trapped during a very long time τ and accumulate [Fig. 4(b)]. The momentum distribution can then get a width δp which is found to decrease indefinitely when the interaction

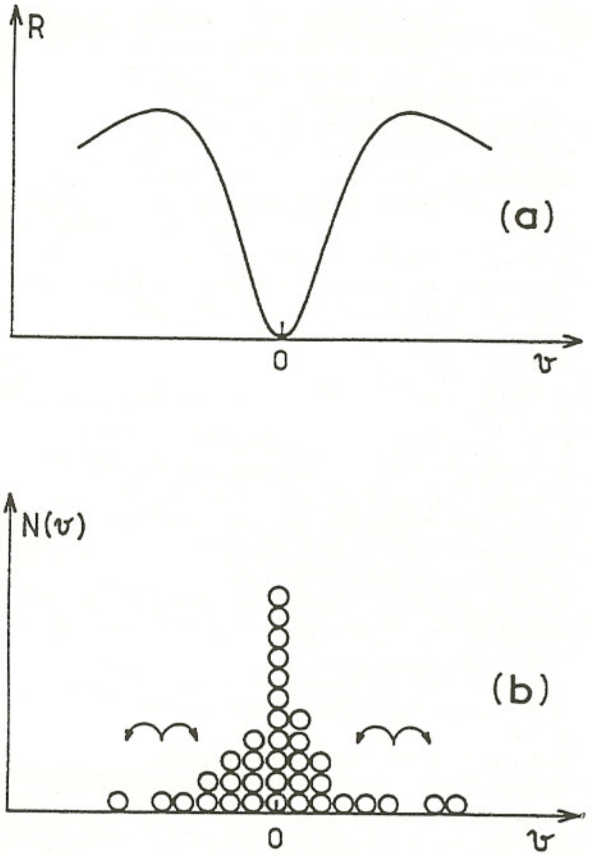

Fig. 4. Principle of subrecoil cooling. a – The photon absorption rate $R(v)$ is made velocity-dependent and vanishes for $v = 0$. b – During their random walk in v-space, atoms can fall in a small interval near $v = 0$ where they remain trapped and accumulate.

time θ tends to infinity. Up to now, two subrecoil cooling schemes have been proposed and demonstrated. In the first one, called "velocity selective coherent population trapping" (VSCPT), the vanishing of $R(v)$ for $v = 0$ is achieved by using destructive quantum interference between different absorption amplitudes [61]. The second one, called Raman cooling, uses appropriate sequences of stimulated Raman and optical pumping pulses for tailoring the appropriate shape of $R(v)$ [62].

4.3.2. *Non ergodic features of the cooling process.* Quantum Monte Carlo simulations using the delay function (see Section 4.2) have provided new physical insight into subrecoil laser cooling [63]. It clearly appears that the random walk of the atom in velocity space is anomalous and dominated by a few rare events whose duration is a significant fraction of the total interaction time. A simple analysis shows that the distribution $P(\tau)$ of the trapping times τ in a small trapping zone near $v = 0$ is a broad distribution, with power-law tails. These tails decrease so slowly that the average value $\langle\tau\rangle$ of τ (or the variance) can diverge. In such cases, the central limit theorem (CLT) can obviously no longer be used for studying the distribution of the total trapping time after N entries in the trapping zone separated by N exits.

It is possible to extend the CLT to broad distributions with power-law tails [64, 65]. We have applied the corresponding statistics, called "Lévy statistics", to subrecoil cooling and shown that one can obtain in this way a better understanding of the physical processes as well as quantitative analytical predictions for the asymptotic properties of the cooled atoms in the limit when the interaction time θ tends to infinity [63, 66]. One important feature revealed by such an analysis is the non ergodicity of the cooling process. As long as the interaction time θ can be, there are always atomic evolution times (trapping times in the small zone around $v = 0$) which can be longer than θ. Another advantage of such a new approach is to allow the parameters of the cooling lasers to be optimized for given experimental conditions. For example, by using different shapes for the laser pulses used in one-dimensional subrecoil Raman cooling, it has been possible to reach for Cesium atoms temperatures as low as 3 nK [67]. The long coherence lengths achieved in this way (several microns) have been very useful for preparing atomic Bloch states in the periodic optical potential produced by a laser standing wave and for observing Bloch oscillations [44] (see Section 3.3.2).

4.3.3. *Recent experimental developments of VSCPT.* The first VSCPT experiments were using an atomic beam of metastable Helium atoms [61]. By starting from a cloud of trapped and precooled atoms, it has been possible to increase by nearly two orders of magnitude the interaction time θ between the atoms and the VSCPT beams, which is important for lowering the temperature (which is predicted to vary as $1/\theta$) and for extending VSCPT to two and three dimensions (see [68] and references therein).

One of the important features of VSCPT is to prepare the cooled atom in linear superpositions of a certain number of coherent wave packets. By using laser induced adiabatic passage, it is possible to coherently manipulate these wave packets [69]. One can for example transfer the whole atomic population into a single atomic wave packet, or two, whose directions can be chosen at will. Very recently, a new method has been developed for directly measuring, from the overlap of two VSCPT wave packets, the spatial correlation function of the cooled atoms [70]. Such a method, which is somewhat analogous to Fourier spectroscopy in optics, leads to a determination of the atomic momentum distribution much more precise than with the usual time of flight techniques. Temperatures as low as 1/800 of the recoil temperature T_R (defined by $k_B T_R/2 = \hbar^2 k^2/2M$) have been measured, and quantitative comparison with the theoretical predictions of Lévy statistics have been made. It has also been possible to recombine two atomic wave packets, after they have flown apart, and to observe interference effects directly demonstrating their coherence.

5. Conclusion

Our ability to manipulate atoms with photons has considerably increased during the last few years and has opened the way to many new applications. One can reasonably expect that further progresses will be achieved thanks to the continuous development of experimental tools, such as better solid state sources, better detectors . . .

At least, two directions of research look promising. First, a better control of "pure" situations involving a small number of atoms and photons in well defined states exhibiting quantum correlations or entanglement. In that perspective, atomic, molecular and optical physics will continue to play an important role by providing a "testing bench" for improving our understanding of quantum phenomena. A second interesting direction is the investigation of new systems, such as Bose condensates involving a macroscopic number of atoms in the same quantum state. One can reasonably hope that new types of coherent atomic sources (sometimes called "atom lasers") will be achieved, opening the way to interesting new possibilities.

It is clear finally that all these developments will strengthen the connections which are being established between atomic physics and other branches of physics, such as condensed matter or statistical physics. The use of Lévy statistics for analyzing subrecoil cooling is an example of such a fruitful dialog. The interdisciplinary character of the present researches on the properties of Bose condensates is also a clear sign of the increase of these exchanges.

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