Phase Modulation of Atomic de Broglie Waves

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Cesium atoms prepared in a state of well-defined total energy have been reflected from a vibrating mirror, causing the matter waves to be phase modulated. The mirror is an evanescent light wave whose intensity is modulated in time at a frequency ν in the range 0 to 2 MHz; the atoms are reflected at normal incidence. The resulting beam consists of a "carrier" plus various sidebands corresponding to de Broglie waves propagating at different velocities. The precision and flexibility of this technique make it a promising tool in atom optics.

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It has been known since the days of de Broglie that the motion of particles of matter is accurately described by a wave equation. Recent experiments in atomic physics have concentrated on the spatial aspect of this wavelike motion, showing that atoms are subject to phenomena such as interference and diffraction by slits and gratings [1,2]. In this Letter we focus rather on the temporal aspect of the wave nature of atomic particles. We present an experiment demonstrating the phase modulation of de Broglie waves by a vibrating potential. Cesium atoms prepared with a well-defined total (kinetic + potential) energy E_i are reflected on an atomic mirror vibrating at frequency ν, and we detect the presence of several energy components $E_f = E_i + nh\nu$ (n positive or negative integer) in the resulting atomic beam.

This device opens up new possibilities for precision experiments in atom optics, in complement to current gratings and slits, since one can take advantage of the fact that frequency and time intervals can be produced more accurately than distance intervals. This feature is now basic to high precision methods in photon optics, such as frequency chains using electro-optics modulators. In neutron optics, direct frequency transfer can be achieved by reflection of a neutron beam off a crystal lattice vibrating at one of its resonant frequencies [3,4]. For atoms or molecules, up to now, experiments using precise rf techniques have relied on the presence of a narrow resonant transition between internal states of the system. By contrast, in our experiment, the internal structure of the atom plays no role in defining the modulating frequency, and the latter can be varied continuously.

Cesium atoms released from a magneto-optical trap (MOT) form a cold "beam," moving vertically downwards. They are velocity selected to have a kinetic energy $E/h = 10.42 \pm 0.03$ MHz at the mirror surface. The mirror potential is produced by the ac Stark shift of the atoms in an evanescent light field propagating along the surface of a glass prism [5]. The amplitude of the light field is modulated at a chosen frequency between 0 and 2 MHz. The presence of several frequency components in the reflected atomic beam is deduced by recording the time taken for the beam to rise and then fall back down to a fixed "probe" height in the Earth's gravitational field. Since the matter waves' group velocity (at any given height) depends on their frequency, the negative-order sidebands arrive first at the probe, followed by the carrier, followed by the positive sidebands. The presence of these sidebands clearly demonstrates the quantum nature of the phenomenon, since they could not appear in the reflection of a classical point-like particle on a modulated potential. They demonstrate also a direct transfer of a precisely synthesized rf frequency onto the atomic motion.

As for any scattering process, the maximum frequency at which efficient transfer can occur between the carrier and the various sidebands is of the order of τ^{-1}, where τ is the interaction time between the atom and the oscillating potential [6]. The potential created by the evanescent wave (EW) varies with height z as $\exp(-2\kappa z)$, where $1/\kappa \simeq 0.19\ \mu$m. The observed atoms fall with mean velocity $v = 25$ cm/s onto this potential, giving $\tau \simeq 1\ \mu$s. Hence, the phase modulation process is only efficient for frequencies up to $\sim$1 MHz. Using $(p + \Delta p)^2/2M = p^2/2M + nh\nu$, this gives momentum transfers Δp of order $4\hbar\kappa$ for $n = 1$. Regarding the phase modulator as a coherent beam splitter, this represents a useful, efficient beam splitting by several times the photon-recoil momentum.

Our vacuum system is based on two glass cubes of side 10 cm, one positioned 70 cm above the other (see Fig. 1). Each is evacuated by a 25 l/s ion pump, and they are connected through a narrow glass tube (ϕ 9 mm, length 140 mm) to allow differential pumping. This system allows us to produce a good vacuum in the lower cell ($\leq 3 \times 10^{-9}$ mbar) while having sufficient vapor pressure of cesium in the upper cell (6×10^{-8} mbar) to load a MOT there in a short time [7].

Light in the experiment is provided wholly by diode lasers at the cesium resonance line (852 nm). The experimental cycle begins by loading 3×10^8 atoms into the MOT in the upper cell during 1 s. The collected atoms are then cooled to 5 μK in optical molasses: the laser intensity is switched to 1 mW/cm^2 per beam, the magnetic field gradient is cut, and the detuning δ between the laser and atomic frequencies is ramped to $\delta = -9\Gamma$, where $\Gamma = 2\pi \times 5.3$ MHz is the FWHM of the atomic

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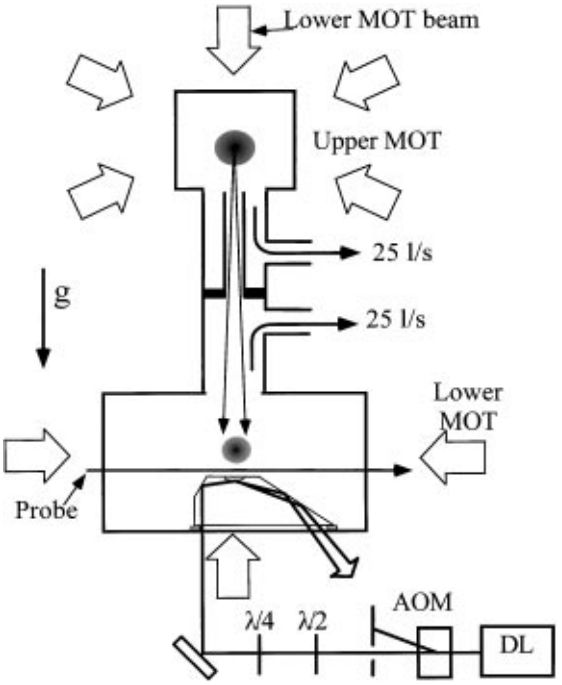

FIG. 1. Experimental setup. Atoms are captured in the upper MOT from a cesium vapor, and transferred to the lower MOT in a high vacuum cell. They are then cooled and dropped into the evanescent wave mirror, which is formed using light produced by a diode laser (DL) and time modulated by an acousto-optic modulator (AOM).

resonance. The trap light is then switched far from the atomic resonance, and the atoms fall into the lower chamber. After a free fall time of 365 ms, a MOT in the lower cell is switched on for 100 ms; it can easily bring to rest the falling atoms, though not all of them are captured, perhaps due to imbalanced light forces as the atoms enter the trapping beams. The overall transfer efficiency is about 20%, so this provides 6×10^7 atoms each 1.4 s, with a lower trap lifetime $\sim$12 s. Once the atoms are caught in the lower MOT, they are compressed to a density 5×10^{10} cm^{-3} by reducing the laser intensity in the MOT to 0.5 mW/cm^2 at $\delta = -3\Gamma$ during 3 ms, and then cooled to 3.6 μK at $\delta = -9\Gamma$ in optical molasses during 20 ms. We used our velocity-selection technique, described below, to measure this temperature. The quadrupole field for the lower MOT is provided by two coils mounted on an *XYZ* translation stage, enabling the trap to be accurately positioned 3.3 mm above the center of the EW atomic mirror. The latter is made by a 100 mW elliptical Gaussian beam, totally internally reflected in a superpolished fused silica prism. The optimum polarization was found to be linear in the plane of reflection. The resulting reflective spot is circular, of radius 400 μm. The prism surface is concave, with radius of curvature 2 cm [8]. The EW intensity is controlled by an acousto-optic modulator (AOM) used in the zeroth order. By sending to the AOM an amplitude-modulated rf signal, the EW intensity $I(t)$ can be made to oscillate at the modulation frequency ν, $I(t) = I_0[1 + \epsilon \cos(2\pi\nu t)]$, with a modulation depth ϵ up to 82%. The AOM is also used to cut the EW—an attenuation to $I_{\min}/I_{\max} = 10\%$ reduces the mirror potential sufficiently so that atoms released from a 3.3 mm height are not reflected.

The atomic de Broglie waves released from the lower MOT have a frequency spread $\Delta E_i/h \sim 7$ MHz, which is greater than the maximum efficient modulation frequency. To obtain well-resolved signals, we therefore energy select the atoms before phase modulating them, using a sequence of two short pulses of the EW intensity (Fig. 2). Each pulse has a duration $\Delta T = 0.4$ ms and their centers are separated by $T = 52$ ms, with the first pulse centered $T/2 = 26$ ms after the atoms have been released from the MOT. The velocity distribution which is thus selected is triangular, centered at $v_i = gT/2 = 25.5$ cm/s, with a width (HWHM) $\Delta v = g\Delta T/4 = 0.098$ cm/s. This corresponds to $\Delta v = 0.28 v_{\rm rec}$, where the recoil velocity is $v_{\rm rec} = \hbar k/M = h/M\lambda = 3.5$ mm/s. To compare with other velocity selection methods based on an atom-light interaction [9], we have checked that this method still works nicely for pulses as short as $\Delta T' = 0.1$ ms, corresponding to $\Delta v' = 0.07 v_{\rm rec} = 250$ μm/s. This finer selection was not useful for the present study since the experimental resolution was then limited by the probe.

During the second pulse, the intensity of the evanescent wave is modulated. The final velocity distribution is probed by a time of flight (TOF) method. We measure as a function of time the absorption of a probe beam which is switched on 32 ms after the second pulse. This probe beam has a truncated Gaussian profile with a vertical width $\delta H = 100$ μm, and it is located at a height $H \simeq 1.7$ mm above the mirror. The time resolution δT of the detection is then given by the transit time of the atoms crossing the probe with a velocity $v_p \sim 18$ cm/s : $\delta T = \delta H/v_p = 0.5$ ms.

The TOF signal shown in Fig. 3(a) was obtained with a nonmodulated mirror ($\epsilon = 0$). It represents the average of 1000 shots, with $\sim$4000 atoms contributing per shot.

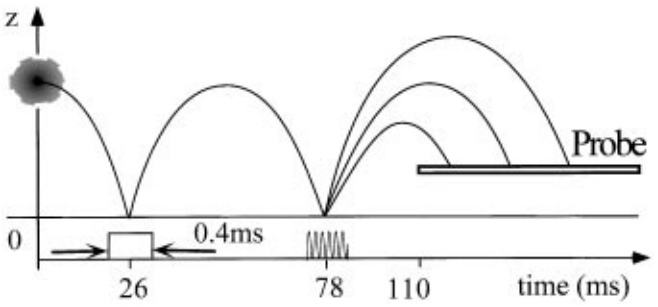

FIG. 2. Atomic trajectories (upper part) and EW intensity (lower part) as a function of time. Atoms with a well-defined energy are selected using a two-pulse technique. The modulation of the EW during the second pulse introduces sidebands onto the de Broglie waves. These are detected by their time of flight to a probe beam introduced after the second pulse.

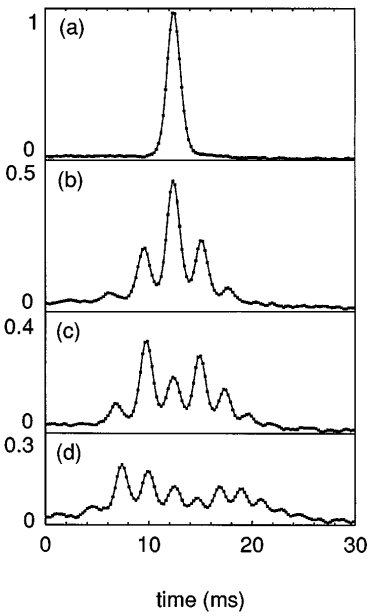

FIG. 3. Time-of-flight signals. (a) Nonmodulated potential; (b)–(d) modulated potential, with frequency $\nu = 950$ kHz (b), 880 kHz (c), and 800 kHz (d). The modulation depth is $\epsilon = 0.82$. The width of the points indicates the time window for digitization and averaging of the continuous signal.

The position of the central feature allows a precise determination of the height of the probe, $H = 1.73$ mm. A Gaussian curve fitted to this feature has a standard deviation $\tau = 0.64$ ms, resulting from the width of the selection pulses and finite resolution of the probe. A small broader feature (height ~2% of the main peak, full width ~ 20 ms) also appears. It may originate from atoms which have been heated during the bouncing process by photon scattering which led to a change of their Zeeman sublevel, and therefore a change of their potential energy in the evanescent field. Atoms close to the mirror when it is turned on or off contribute to an even smaller wider background. This TOF spectrum [3(a)] is used in the following as an instrument function $\phi(t)$ describing the response of our system.

The TOF spectra [3(b)–3(d)] were obtained with various modulation frequencies ν. The average intensity I_0 is the same as in Fig. 3(a), and the modulation depth for 3(b)–3(d) is $\epsilon = 0.82$. We have fitted these TOF signals by adding, with adjustable position and weight, several instrument functions $\sum_n a_n \phi(t - t_n)$. We obtain from this fit the positions t_n of the sidebands, from which we derive the corresponding energy transfers. These agree with the theoretical predictions [Fig. 4(a)]. From our measurements, we deduce $h/Mg^2 = (3.14 \pm 0.02) \times 10^{-11}$ s^3, to be compared with the expected value 3.120×10^{-11} s^3, using the local $g = 9.81$ m/s^2 in Paris.

The weight of the sidebands is also in reasonable agreement with the theoretical predictions. If we approximate the diffraction by the modulated mirror as diffraction from a thin phase grating [6], we get for the weight of the nth band $|a_n|^2 = |J_n(u)|^2$, where $u = \epsilon\beta(2\pi\nu/\kappa v_i)Mv_i/\hbar\kappa$, and $\beta(x) = (\pi x/2)/\sinh(\pi x/2)$. This prediction is plotted in dotted lines for $n = 0, \pm 1$, together with the experimental results, in Figs. 4(b) and 4(c). The discrepancies appear mainly at high frequency where the energy transfer is more efficient than predicted. This increased efficiency can be recovered theoretically from a numerical integration of the Schrödinger equation, taking into account the van der Waals interaction and the abrupt termination of the potential at the glass surface. This introduces shorter time scales, and therefore increased bandwidth, into the system.

Two effects cause asymmetries in the spectra [Figs. 3(b)–3(d)]. These are the nonlinear dispersion relation for de Broglie waves [6] and the loss of atoms falling at the edge of the Gaussian spot. The latter reduces mostly the contribution of atoms with upshifted energy, as can be shown by solving the classical equations of motion in the three-dimensional mirror potential.

These experiments have been performed with a detuning of the EW equal to $\delta/2\pi = 1$ GHz, which is sufficient to ensure that spontaneous emission during the bouncing process plays a minor role. The probability per bounce for such a process is $P = \Gamma M v_i/\hbar\kappa\delta \simeq 0.5$ [8,10], and the

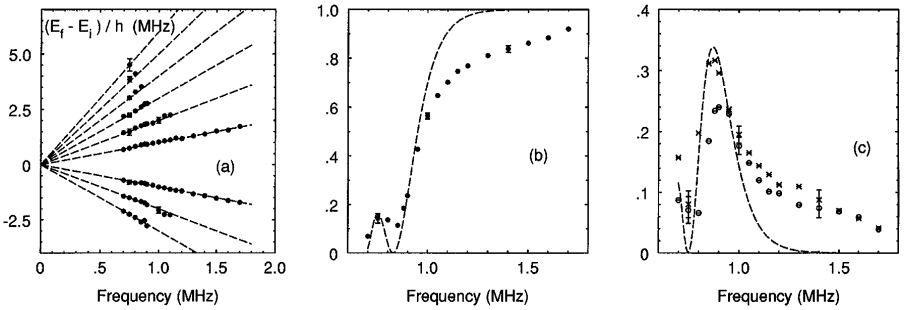

FIG. 4. (a) Energy transfers (deduced from arrival times at the probe) as a function of modulation frequency. The dotted lines show the theoretical prediction $\Delta E_n/h = n\nu$, with no fitted parameters. (b) Amplitude of the carrier and (c) amplitude of the sidebands $n = -1$ (×) and $n = 1$ (○), as a function of modulation frequency; the dotted curves show an approximate theoretical prediction for reflection off an oscillating pure-exponential potential. The more efficient transfer seen experimentally above 1 MHz is related to the cutoff of the light shift potential at the dielectric surface. Experimental error bars, obtained by an analysis of the fitting procedure described in the text, have been plotted in (a)–(c) for all points at $\nu = 0.75$, 1, and 1.4 MHz.

corresponding random recoil is small compared to the velocity separation of two adjacent sidebands. In addition, we estimate the probability for photon scattering due to the stray light scattered at the total internal reflection to be lower than 0.05 per atom during the whole time sequence. The probability of spontaneous emission processes could be much reduced if dielectric coatings or surface plasmons were used to enhance the evanescent wave intensity, therefore allowing a larger detuning [11–13].

This phase modulator for de Broglie waves is the temporal equivalent of a spatial periodic phase grating. This first component for atom optics in the time domain can be complemented by other devices working along similar lines, making temporal equivalents of lenses or prisms. They are obtained by moving the mirror surface (i.e., by varying the intensity of the EW) in order to adjust as a function of time the momentum transferred to the atom [14]. A thin sheet of light with a time modulated intensity crossing an atomic beam can also provide the equivalent for de Broglie waves of an electro-optic modulator [15]. If needed, all these time-domain components can be made selective with the atomic internal state, for a proper choice of light polarization and detuning. They do not require any fabricated material mask, and they have a transmission factor of unity.

Another prospect deals with the two-pulse selection method presented above. It is a convenient way to select a chosen small region in the atom's phase space, and it can be used to study the phase space distribution of the MOT source, or to prepare diffraction limited beams. For instance, using pulses of width 0.1 ms, we have obtained a beam having a flux ~500 atoms per second with $M\Delta z\Delta v \simeq 6\hbar$, where $\Delta z = v_i \Delta T/2$ is the HWHM of the vertical atomic distribution just after the bounce. The fundamental limitation of this phase space selection method ($M\Delta z\Delta v > \hbar/2$) originates from the diffraction in time which appears when the pulse width is very small. This phenomenon of diffraction by a time slit has been discussed in detail in the context of neutron optics [16–18]. To our knowledge, it has not been directly observed for material particles.

Finally, this phase modulator can be used for new atomic interferometers [6,19]. For instance, two interfering paths could involve N and $N + 1$ bounces, with the initial splitting and final recombination carried out by means of modulated pulses. The interference phase is then sensitive to the acceleration due to gravity, and this could be used to measure g precisely [20]. Alternatively, the well known three-grating interferometer [21,22] can be transposed into the time domain using 3 modulated pulses to split, recombine, and analyze the wave function. Because of its symmetry, this design is insensitive to the initial velocity dispersion and to g. It would be a precise tool to measure the phase shifts during the bouncing process itself [23], including the contribution of the van der Waals interaction between the atom and the dielectric.

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