Optical-Sideband Cooling of Visible Atom Cloud Confined in Parabolic Well

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An assemblage of < 50 Ba^+ ions, contained in a parabolic well, has been visually observed and cooled by means of near-resonant laser irradiation.

Taking the attitude that the pursuit of as basic an ideal as a single atomic particle at rest in space[1] is a thoroughly worthwhile intellectual endeavor we are undertaking experiments along these lines. We hope to make a single (charged) atom trapped in a parabolic well visible[2] and simulataneously freeze out its motion in the well[3] by means of suitable laser beams. A stationary atom has obvious advantages in the suppression of transit-time broadening and of first- and second-order Doppler shifts[1,3] which are limiting high-resolution spectroscopy and its application to the measurement of time. Radiation-pressure cooling of an ordinary atomic gas and, independently, the (optical) sideband cooling of a single ion in a parabolic well[3] have been proposed[4] and/or demonstrated[5] before.

The ions were confined in a miniaturized Paul rf trap[1] (see Fig. 1) with a cap-cap separation $2Z_0 \simeq 0.5$ mm. The trap, Ba oven, and electron gun were housed in a Pyrex envelope exhausted by an 8-L VacIon pump. For an applied trapping voltage of amplitude $V_0 \simeq 200$ V at $\Omega = 2\pi \times 18$ MHz an axial oscillation frequency in the well $\omega_w \simeq 2\pi \times 2.4$ MHz was measured via ion ejection caused by a resonant ω_w drive. For the purpose of filling the well a barium atomic beam and a 900-V electron beam, both entering the trap through the gap between ring and cap, were crossed in its center. Initially the ions were thermalized through viscous-drag cooling[1] provided by a light buffer gas at $\lesssim 10^{-4}$ Torr. To this end the pump was turned off and He gas was admitted through a He-Pyrex diffusion leak heated to dark red heat. The ions were illuminated with the principal laser beam of wavelength $\lambda_0 = 493.4$ nm and an auxiliary beam at $\lambda_a = 649.9$ nm focused along the $-z'$ direction into the well center, with waists $w \simeq 60$ μm and powers ~ 2 and 5 mW. cw dye lasers of spectral width ~ 100 MHz and ~ 100 GHz using coumarine-102 and rhodamine-B, respectively, provided the λ_0 and λ_a beams. A discharge tube containing Ba vapor and ~ 10 Torr helium served as a frequency standard for the Ba^+ $6^2S_{1/2}$-$6^2P_{1/2}$ and $5^2D_{3/2}$-$6^2P_{1/2}$ transitions. Both λ_0 and λ_a were continuously adjusted for maximum fluorescence observed in the standard. On the cloud in the well we monitored the blue resonance fluorescence. With the beam oven in equilibrium it took about 20 sec for the cloud of ~ 50 μm diameter to reach its maximum brightness. The cloud was then easily observed through a microscope visually (see Fig. 2).

Now, for the optical cooling experiment, the He leak was turned off and the ion getter pump operating at 9 kV was turned on while the cloud was continually watched visually. As the buffer gas pressure fell, the fluorescence weakened and finally disappeared, even though both Ba oven and electron beam were on. The pump current (leakage dominated) then indicated a pressure < 10^{-8} Torr. This suggested very short dwell times τ_F in the focal region for the now unbuffered ions, presumably as a result of heating processes[1] always present in rf traps. However, when the blue laser was detuned to lower frequencies, $\omega = \omega_0 - \Delta_{1/2}$, until the fluorescence in the standard

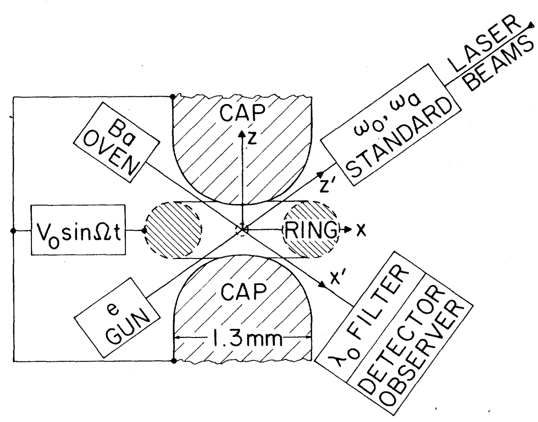

FIG. 1. Apparatus for trapping, cooling, and visually observing Ba^+ cloud (schematic). Ba oven and electron gun were actually not in the x-z plane.

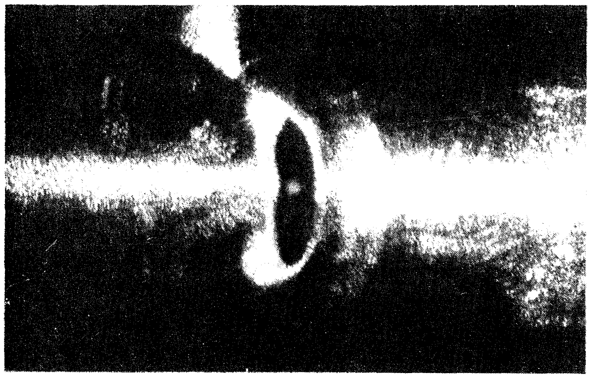

FIG. 2. Photograph of trap looking in $-x'$ direction. The Ba^+ cloud is visible in the center.

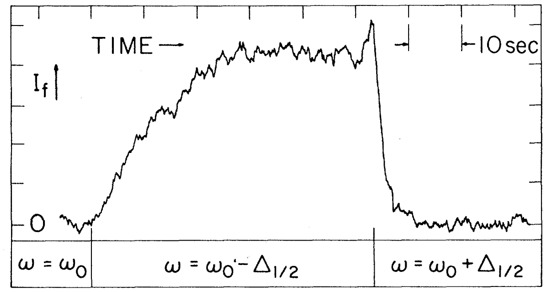

FIG. 3. Ba^+ fluorescence intensity I_f vs time. The trap, empty when the laser frequency ω coincides with that of the resonance line ω_0, fills slowly for $\omega = \omega_0 - \Delta_{1/2}$ and empties rapidly for $\omega = \omega_0 + \Delta_{1/2}$.

decreased to about one half of its maximum value the trap again began to fill, but only with the red λ_a laser on. Presumably the laser irradiation greatly increased τ_F through cooling. This was corroborated when the trap was filled with the lasers at $\omega_0 - \Delta_{1/2}$ and ω_a and later the oven and electron gun were turned off. For $\Omega = 2\pi \times 58$ MHz, $\tau_F > 5000$ sec was observed then photometrically. The λ_a beam merely served to clean ions out of the $D_{3/2}$ level into which they fall one-third of the time from the excited $P_{1/2}$ level. When this beam is blocked, the ions quickly collect in the metastable $D_{3/2}$ level and get lost from the trap as a result of various heating processes which now are no longer compensated for by the optical cooling. Fast chopping of the red light at ~100 Hz caused no ion loss and provided an "atomic" modulation of the blue fluorescence signal for use with a lockin detector. The trace of the scattered blue-light intensity versus time (see Fig. 3) was obtained in this way. The graph shows slow filling of the trap with the blue laser at $\omega_0 - \Delta_{1/2}$ and the fast emptying expected when it was switched to $\omega_0 + \Delta_{1/2}$ and the laser beam heated the ions. In subsequent runs fluorescence disks *as small as < 5 μm due to a single ion* were also seen but will be discussed in another paper.

Under our experimental conditions an ion in the rf trap sees essentially a three-dimensional ellipsoidal potential well[1]

$$\psi(x,y,z) = D(x^2+y^2+4z^2)/4Z_0^2,$$

with $D = 9.4$ V. Actually the z motion of the ions is given by $z = \bar{z}(t) + \zeta(t)$. It is composed of the free secular motion of the guiding center, $\bar{z}(t) = z_0 \cos\omega_w t$, and the *forced* micromotion $\zeta(t) = \zeta_0 \cos\omega_w t \cos\Omega t$ with $\zeta_0 \equiv \sqrt{2}\, z_0\omega_w/\Omega \ll z_0$. The basic idea of sideband cooling is simple. To an atom oscillating in a parabolic well at ω_w, the spectrum of a monochromatic incident wave of frequency ω through the Doppler effect appears to have components at $\omega + n\omega_w$, $|n| = 0,1,2,\ldots$. When one of the sidebands $\omega + n\omega_w \equiv \omega_n'$ falls near the resonant frequency ω_0 of the atom in its rest frame the transition is excited and $\hbar\omega$ photons are absorbed. The atom reemits photons of *average* energy $\hbar\omega_n'$ the reemission spectrum components at $\omega_n' \pm m\omega_w$ being symmetric about ω_n', $m = 0,1,2,3,\ldots$. For $n > 0$ the average energy defect per scattered $\hbar\omega$ photon $\hbar(\omega_n' - \omega) = \hbar n\omega_w$, must be provided by the free oscillation in the well which is thereby cooled. Analogous heating occurs for $n < 0$. Restricting ourselves to a one-dimensional model, we see that the frequency ω of a plane electromagnetic wave moving along the positive z axis because of the Doppler effect will appear to a moving ion as shifted to

$$\omega' = \omega - \dot{z}/\lambda\!\!\!^{-}, \quad \lambda\!\!\!^{-}\omega = c.$$

The corresponding phase is given by $\varphi = \int \omega' dt = \omega t - z(t)/\lambda\!\!\!^{-}$. For the electric field seen by the atom in its own rest frame one has $E_y = \cos\varphi$. Comparing displacement $z(t)$ with the characteristic length $\lambda\!\!\!^{-}_0 \simeq \lambda\!\!\!^{-}$, we now distinguish two regimes: the Lamb-Dicke regime, $|z| \ll \lambda\!\!\!^{-}_0$, and the Doppler regime, $|z| \gg \lambda\!\!\!^{-}_0$. In the Lamb-Dicke regime $\cos\varphi$ is easily expanded:

$$E_y \simeq \cos\omega t - (z/\lambda\!\!\!^{-}_0)\sin\omega t,$$

showing the dominance of the unshifted carrier at ω over the sideband background, the total relative power content $|z/\lambda\!\!\!^{-}_0|^2$ of which may be made arbitrarily small. An interesting consequence of this important equation is that the spectrum of an atom, though moving around on an arbitrary path at an arbitrarily high speed, will show no appreciable Doppler effect, provided it only remains in a region much smaller than $\lambda\!\!\!^{-}_0$. This vanishing of the first-order Doppler effect has been discussed or observed in high-resolution spectroscopy with neutron waves on nuclei and with electromagnetic waves on atoms and nuclei in the microwave, radio, and γ-ray regions since 1939.

Returing to our ion motion in the well we find $\omega' \simeq \omega - \dot{\bar{z}}/\lambda\!\!\!^{-}_0 - \dot{\xi}/\lambda\!\!\!^{-}_0$. In very good approximation the expression for ω' may be interpreted as describing an experiment in which an equivalent $\omega_e(t)$ wave, frequency modulated by the forced ξ motion, with $\omega_e \equiv \omega - \dot{\xi}/\lambda\!\!\!^{-}_0$, is incident upon an atom undergoing a free purely harmonic translational oscillation at ω_w. In the Lamb-Dicke regime we have $z_0/\lambda\!\!\!^{-}_0 = \beta \ll 1$ and $\xi/\lambda\!\!\!^{-}_0 \ll \beta \ll 1$. Again the frequency modulation, now due to the micromotion, has no effect on the spectrum of the original wave of frequency ω. This implies, for the field seen by the atom, $E_y \simeq \cos[\omega t - \beta \cos\omega_w t]$. The usual expansion of this yields spectral components at $\omega \pm n\omega_w$ of relative power $J_n^2(\beta)$, with $J_n(\beta) \to 0$ for $\beta < n$. For $\beta \ll 1$, the atom of resonance frequency ω_0 sees spectral components at ω and $\omega \pm \omega_w$ of relative powers 1 and $f_\pm$ only, which will independently excite the $\Delta\omega_0$-wide atomic line of shape $g(\omega) = \{1 + [2(\omega - \omega_0)/\Delta\omega_0]^2\}^{-1}$ at rates Sg_ω, f_+Sg_+, and f_-Sg_-, with $g_\omega \equiv g(\omega)$ and $g_\pm \equiv g(\omega \pm \omega_w)$. A laser width $\Delta\omega_L \ll \Delta\omega_0$ is assumed here. As quantum effects are important, we express $f_\pm$ in terms of the oscillator quantum number v. For the oscillation energy we have $W = (M/2)\omega_w^2\beta^2\lambda\!\!\!^{-}_0{}^2 = (v + 1/2)\hbar\omega_w$, and $f_\pm = \beta^2/4 = (v + 1/2)/v_L$ follows, with $v_L \equiv 2(\lambda\!\!\!^{-}_0/\lambda\!\!\!^{-}_C)(\omega_w/\omega_0)$; here $\lambda\!\!\!^{-}_C$ denotes the Compton wavelength of the atom. This requires modification. During the Doppler-effect-induced $\omega \pm \omega_w$ excitation, v changes to $v \mp 1$ which is accompanied by a corresponding change in the sideband powers $f_\pm$ while the transition takes place. In order to allow for this we use $(v_{\rm final} + v_{\rm initial})/2$ in place of v in the formula for $f_\pm$ and obtain $f_\pm = (v + \frac{1}{2} \mp \frac{1}{2})/v_L$. Similarly, for excitation at ω, during reemission ω, $\omega + \omega_w$, and $\omega - \omega_w$ photons leave the atom at relative rates $1 + 1$, f_+, and f_-; one of the two 1's reflects the unshifted emission in the x direction. This entails that on the average the oscillatory motion retains the additional energy $\hbar\omega_w(f_- - f_+)/2$ per scattered photon. The net cooling power is now given by $P_c \simeq \omega_w S[vg_+ - (v+1)g_- - g_\omega/2]/v_L$. This vanishes for $vg_+ - (v+1)g_- - g_\omega/2 = 0$, yielding $v = v_{\rm min} = (g_w/2 + g_-)/(g_+ - g_-)$. To illustrate we take $\omega + \omega_w = \omega_0$, $\omega_w \gg \Delta\omega_0$, and get $g_+ = 1$, $4g_\omega \simeq (\Delta\omega_0/\omega_w)^2$, $g_- \simeq g_\omega/4$, $v_{\rm min} \simeq \frac{3}{16}(\Delta\omega_0/\omega_w)^2 \ll 1$, and a minimal average energy

$$W_{\rm min} \simeq 3\hbar\Delta\omega_0(\Delta\omega_0/\omega_w)/16 \ll \hbar\Delta\omega_0$$

above the zero-point energy. For our *current* experiment we have $\Delta\omega_0 \gg \omega_w$, and the choice $\omega_0 - \omega = \Delta\omega_0/2$ seems desirable. This yields $g_\omega = \frac{1}{2}$, $g_- \simeq \frac{1}{2}$, $g_+ - g_- \simeq 2\omega_w/\Delta\omega_0$, and finally $v_{\rm min} \simeq \Delta\omega_0/2\omega_w$ and $W_{\rm min} \simeq \hbar\Delta\omega_0/2$. Also $\beta_{\rm min}^2 \simeq (\Delta\omega_0\omega_0/\omega_w^2)(\lambda\!\!\!^{-}_C/\lambda\!\!\!^{-}_0)$ is of interest. With the numerical values $\Delta\omega_0/2\pi = 15$ MHz, $\omega_w/2\pi \simeq 2.5$ MHz, $\lambda\!\!\!^{-}_0 \simeq 5\times10^{-5}$ cm, $\lambda\!\!\!^{-}_C \simeq 10^{-15}$ cm, $\omega_0/2\pi \simeq 6\times10^{14}$ Hz, $v_L \simeq 400$, we have $v_{\rm min} \simeq 3$, $\beta_{\rm min} \simeq 0.17$, showing that indeed we may be able to operate in the desirable Lamb-Dicke regime soon. Contenting ourselves with a laser power P^* simulating a thermal radiation field at temperature $T^* = \hbar\omega_0/k$, we realize a fractional upper-level population $(1+e)^{-1}$ and a scattering rate $S = S^* \simeq \Delta\omega_0/(1+e)$. For large v and $S = S^*$ a characteristic sideband cooling time τ_c^* may be obtained, $1/\tau_c^* \simeq [(g_+ - g_-)/v_L(1+e)]\Delta\omega_0$. For P^* one finds $P^* \simeq (w/\lambda\!\!\!^{-}_0)^2 \times h\omega_0\Delta\omega_0$ where w is the beam diameter at the site of the ion. Our parameters give numerical values $\tau_c^* \simeq 10^{-4}$ sec and $P^* \simeq 0.1$ μW.

So far we have experimental data only for $\omega_0 - \omega = \Delta_{1/2} \simeq 2\pi \times 500$ MHz. Here we should have $v_{\rm min} \simeq \frac{3}{8}(\omega_0 - \omega)/\omega_w = 75$ and $2z_0 \simeq 4\lambda\!\!\!^{-}_0(v_{\rm min}/v_L)^{1/2} \simeq 0.14$ μm, probably realized in our one-ion runs already. The complex multi-ion mode where we see $2z_0 \simeq 50$ μm seems dominated by rf ion-ion heating.[1] Here we offer these estimates. Resonance occurs for $\omega' = \omega_0$ or, with $\omega_{0e} \equiv \omega_0 + \dot{\bar{z}}/\lambda\!\!\!^{-}_0$, for $\omega_e(t) = \omega_{0e}(t)$. This may be interpreted as the description of the excitation of a resonance at $\omega_{0e}(t)$ by a wave of frequency $\omega_e(t)$. For $\Delta\omega_0 \approx \Omega \gg \omega_w$ we now approximate the $\omega_e(t)$ spectrum roughly by a quasithermal $2\beta\omega_w$-wide band at T^* centered at ω. Appreciable cooling begins when $\omega_{0e}(t)$ penetrates into this band for $\beta > \beta_a \equiv (\omega_0 - \omega)/2\omega_w \simeq 100$, $2z_0 > 16$ μm, and $W > 10$ meV. Cooling power P_c increases with β, goes through a maximum for $\beta = 2\beta_a$, and then falls off. For $\beta = 2\beta_a$ the scattering rate is $S^*/2$, and the average effective energy defect between irradiating and reemitted photons is $\sim\hbar[\omega_0 - \omega]/2$, giving $P_{c\,\rm max} \simeq 0.07\hbar(\omega_0 - \omega)\Delta\omega_0 \simeq 14$ eV/sec. In the expression for P^* we must substitute the width $2(\omega_0 - \omega)$ of the irradiating ω_e spectrum for $\Delta\omega_0$ here and obtain $P^* \simeq 6$ μW.

We estimate < 50 ions in the cloud.[1] At these densities thermalization[1] due to Coulomb collisions takes place in ~ 1 msec. The single laser beam therefore cools all three degrees of freedom rapidly. For a single ion, however, it will be necessary to make all three principal axes of the ellipsoidal potential well different and direct the laser beam more or less along $\hat{i} + \hat{j} + \hat{k}$.

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Quantum-Mechanical Inversion of the Differential Cross Section: Determination of the He-Ne Potential

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The He-Ne potential is obtained from the measured differential cross section by a systematic inversion procedure based on the generalized opitcal theorem. The potential, obtained in the range from 2.1 to 6 Å, is in good agreement with recently proposed model potentials. The results constitute the first example of a fully direct and systematic inversion of experimental scattering data of real systems.

Molecular-beam scattering experiments are the source of some of the most accurate and detailed information on intermolecular potentials.[1] In most cases the extraction of the interaction from the data is done by assuming an explicit functional form for the potential that includes free parameters. These quantities are then varied in a trial-and-error procedure until the cross section calculated from the potential yields a good fit to the experimental data.[2] It is, of course, very desirable to replace this method by a direct, systematic inversion procedure. It was found by Buck[3] and Buck and Pauly[4] that part of this task could be achieved by applying the Firsov[5] semiclassical inversion, which generates the potential from the deflection function (or phase shifts). The deflection function itself is, however, obtained from the cross section by a fitting procedure using a presupposed functional form. The purpose of the present article is to demonstrate that a *complete* inversion of scattering data is possible for real systems. A key step in this procedure is the determination of the scattering amplitude, hence the phase shifts, from the experimental cross section by the unitarity method.[6-8] Though the method was tested for simulated and idealized input data, the *practical* application was hampered by the incomplete and inaccurate data available. The advent of a new generation of molecular-beam machines, however, made it possible to obtain very precise data with high angular and velocity resolution.[9] The measured differential cross section for the system He + Ne taken at the energy of $E = 64.4$ meV ($v = 1927$ m/s) and transfomed to the center-of-mass system is shown in Fig. 1(a). This system was chosen be-

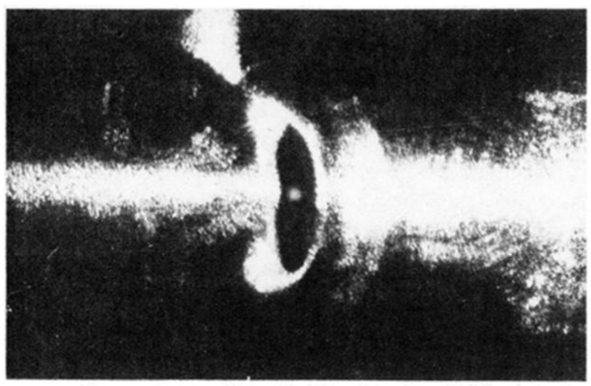

FIG. 2. Photograph of trap looking in $-x'$ direction. The Ba^+ cloud is visible in the center.