Optical Bistability at Low Light Level due to Collective Atomic Recoil

M. Vengalattore,* M. Hafezi, M. D. Lukin, and M. Prentiss

MIT-Harvard Center for Ultracold Atoms, Harvard University, Cambridge, Massachusetts 02138, USA

(Received 14 April 2008; published 8 August 2008)

We demonstrate optical nonlinearities due to the interaction of weak optical fields with the collective motion of a strongly dispersive ultracold gas. The combination of a recoil-induced resonance in the high gain regime and optical waveguiding within the dispersive medium enables us to achieve a collective atomic cooperativity of 275 ± 50 even in the absence of a cavity. As a result, we observe optical bistability at input powers as low as 20 pW. The present scheme allows for dynamic optical control of the dispersive properties of the ultracold gas using very weak pulses of light. The experimental observations are in good agreement with a theoretical model.

DOI: 10.1103/PhysRevLett.101.063901 PACS numbers: 42.65.Pc, 37.10.Vz, 42.50.Gy

Motivated by potential applications to quantum information science [1–4], there have been intense experimental efforts to realize strong nonlinear interactions between dilute atomic ensembles and weak optical fields. Methods to achieve such quantum nonlinear couplings in dissipation-free media have relied mainly on two approaches. First, the interaction time and the coupling between the atoms and the photons can be enhanced by placing the atoms within a high finesse cavity [5,6], an approach that comes at the expense of considerable experimental complexity and low bandwidth. Alternatively, near-resonant light propagating through optically dense media can also result in strong nonlinear interactions. In order to limit the dominant linear absorption, this latter approach requires the use of a coherent multiphoton process such as electromagnetically induced transparency [7]. An added benefit of such a multiphoton process is the enhancement of the interaction time due to slow group velocities. However, since electromagnetically induced transparency relies on quantum interference between the internal states of an atom, it is fairly sensitive to inhomogenous optical and magnetic fields.

In this Letter, we demonstrate nonlinear optical effects due to the interaction between weak pulses of light and the collective motion of an ultracold slow-light medium. Using the motional degrees of freedom to create a highly dispersive gas alleviates the sensitivity to external fields. As shown in recent studies [8,9], the implementation of a recoil-induced resonance (RIR) in an optically dense anisotropic gas allows for strong atom-light interaction due to the combination of slow group velocities and transverse confinement of the optical fields within the atomic medium. Because of this strong coupling, we observe optical bistability at input powers as low as 20 pW, an upper bound limited mainly by the photodetector efficiency.

The motion of delocalized atoms under the influence of two or more light fields can mediate the conversion of atomic kinetic energy into radiation [10]. These processes, termed "recoil-induced resonances," can be described in terms of stimulated Raman transitions between different momentum classes of the atomic ensemble [11]. For an optically thin medium, the atom-light interaction has little effect on the momentum distribution of the gas and a perturbative analysis reveals that the probe field experiences weak absorption (gain) for positive (negative) detuning relative to the pump field.

In contrast, as the optical depth of the atomic gas increases, the strong atom-light coupling leads to large optical gain [8] and significant effects on the momentum distribution of the atomic gas. The strong amplification of the probe field and its subsequent backaction on the collective motional states of the gain medium result in a nonlinear optical response even for weak incident beams. In this work, this backaction is evidenced by the observation of optical bistability in the probe transmission.

The experimental scheme to create a strongly dispersive ultracold gas using the RIR is described in previous work [8]. About 5×10^8 ^{87}Rb atoms are confined in a highly anisotropic magneto-optic trap at typical temperatures of $\sim 20\ \mu$K. In this trap, the atom cloud assumes the shape of a cylinder with an approximate radial (longitudinal) extent of 200 μm (3 cm). The pump and probe beams for the RIR share the same linear polarization and are directed along the long axis of this cylinder to take advantage of the large optical depth (OD) along this axis. The Rabi frequency of the pump beam was typically $\Omega_1/\Gamma = 1.5$, where Γ is the natural linewidth of this transition. At low OD, the transmission spectrum of the RIR exhibits a characteristic dispersion-shaped spectrum with gain (absorption) for detuning $\delta < 0$ ($\delta > 0$). At higher OD, the probe is almost completely extinguished on the absorption side of the resonance, leaving the gain peak as the only distinguishable feature (Fig. 1).

For an atomic gas with high OD, Fig. 2(a) shows the probe transmission as the pump-probe detuning is scanned across the RIR. Depending on the sign of this detuning chirp, there is a shift in the resonance as well as in the maximum gain. This hysteretic nature of the transmission diminishes as the OD is lowered to less than ~ 10. At larger cooperativity, the atomic ensemble and the light fields form a strongly coupled system and the probe interacts with an atomic ensemble whose motional coherences and momen-

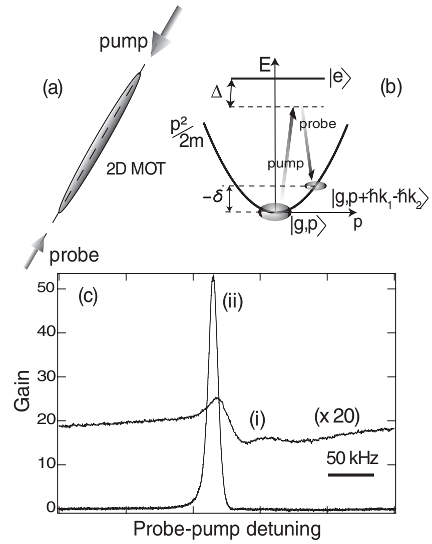

FIG. 1. (a) Schematic indicating the pump and probe beams used for the RIR, where MOT is the magneto-optic trap. (b) The energy levels relevant to the RIR. (c) Absorption spectrum of a probe beam around the RIR for longitudinal optical depth of ~10 (i) and ~40 (ii).

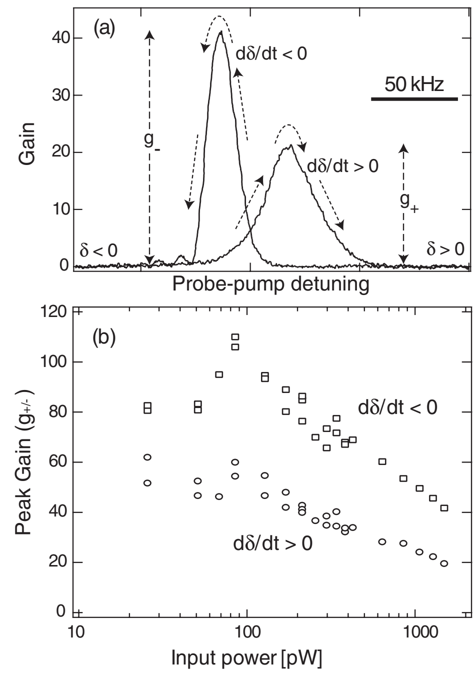

FIG. 2. (a) Absorption bistability due to the interaction between the probe and the collective motion of the atomic gas. Depending on the sign of the detuning chirp ($d\delta/dt$), the transmission indicates a shift in both the resonance and the peak amplification (g_+, g_-). The dashed arrows indicate the direction in which the detuning is scanned across the RIR. (b) Peak probe amplification indicates a bistable response down to the detection limit of the input probe power.

tum distribution are the cumulative result of prior interactions with the light fields. As seen in Fig. 2(b), the bistable response ($g_-/g_+ > 1$) persists down to the detection limit (~20 pW) of the input probe power.

In order to understand this behavior, we first note that the amplification of the probe is accompanied by the transfer of atoms from a momentum p to a momentum $p + 2\hbar k$ [Fig. 1(b)]. Thus, as the detuning is scanned across the gain side of the RIR, atoms at various momenta are brought into resonance with the light fields and transferred to higher momentum states. Crucially, in the case of a negative chirp ($d\delta/dt < 0$), these transferred atoms are brought *closer* to resonance with the light fields as the detuning is scanned. Thus, in this process, atoms are progressively swept up the momentum ladder due to the time-varying detuning, resulting in both an enhanced amplification and a shift in the location of the RIR. In contrast, for a positive chirp ($d\delta/dt > 0$), the atoms are transferred to states that are *farther* from resonance and the probe transmission resembles that obtained for a static thermal distribution.

To quantitatively explain these observations, we use a theoretical model that describes two classical light fields that are coupled to the motional degrees of freedom of an elongated ensemble of two-level atoms. To mimic the experiment, the atoms are assumed to be tightly confined in the radial dimension. Accordingly, only the atomic momentum along the long axis is relevant and the Hamiltonian can be written as [12]

$$\mathcal{H} = \sum_k \bigg[\frac{\hbar^2 k^2}{2m} c_g(k)^\dagger c_g(k) + \bigg(\frac{\hbar^2 k^2}{2m} + \hbar\omega_0 \bigg) c_e(k)^\dagger c_e(k) + i\hbar \sum_{j=1,2} (g_j a_j^* e^{i\omega_j t} c_g(k-k_j)^\dagger c_e(k) - \text{H.c.}) \bigg], \quad (1)$$

where $c_g(k)$ [$c_e(k)$] are the annihilation operators of ground [excited] state atoms with momentum $\hbar k$, ω_0 is the transition frequency of the two-level atoms, g_1 (g_2) is the atom-light coupling coefficient (single-photon Rabi frequency), and a_1 (a_2) corresponds to the normalized electric field of the pump (probe) beams. In the far-detuned limit, the excited states can be adiabatically eliminated. Considering counterpropagating pump and probe fields, we introduce the dimensionless momentum $p = \hbar k/(2\hbar k_0)$ where $2k_0 = k_1 - k_2$. Accordingly, the equation for the coherences and populations of the different momentum classes in the ground state can be written as

$$\begin{aligned} \frac{d}{dt}\rho(p,p') &= 4i\omega_r(p'^2 - p^2)\rho(p,p') + i\frac{g_1 g_2 a_1^*}{\Delta} \\ &\quad \times a_2 e^{-i\delta t}[\rho(p+1,p') - \rho(p,p'-1)] \\ &\quad + i\frac{g_1 g_2 a_1}{\Delta} \\ &\quad \times a_2^* e^{+i\delta t}[-\rho(p,p'+1) + \rho(p-1,p')], \end{aligned} \quad (2)$$

where $\rho(p,p') \equiv \langle c_g^\dagger(k')c_g(k)\rangle$, ω_r is the recoil frequency, and $\delta = \omega_2 - \omega_1$ is the detuning. Also, using the slowly varying envelope and single-mode approximations, the evolution of the probe amplitude can be written as

$$\frac{d}{dt}a_2 = iN\frac{g_1 g_2}{\Delta}a_1 e^{i\delta t}\sum_p \rho(p-1,p) - \frac{\kappa}{2}(a_2 - a_{in}). \tag{3}$$

Retaining only first-order coherence terms between momentum classes, the above equations can be written as a set of coupled equations for the population $\Pi_p = \rho(p,p)$, the first-order coherence $\eta_p \equiv \rho(p+1,p)e^{-i\delta t}$, and the probe amplitude a_2:

$$\begin{aligned}\dot{\Pi}_p &= [-i\beta^* a_2(-\eta_p + \eta_{p-1}) + \text{c.c.}] \\ &\quad - \gamma_{\text{pop}}(\Pi_p - \Pi_{\text{th},p}), \\ \dot{\eta}_p &= i\{4\omega_r[p^2 - (p+1)^2] - \delta(t) + i\gamma_{\text{coh}}\}\eta_p \\ &\quad - i\beta a_2^*(\Pi_{p+1} - \Pi_p), \\ \dot{a}_2 &= i\beta N\sum_p \eta^*_{p-1} - \frac{\kappa}{2}(a_2 - a_{\text{in}}),\end{aligned} \tag{4}$$

where γ_{pop} (γ_{coh}) are the population (coherence) relaxation rates, respectively, and $\beta = g_1 g_2 a_1/\Delta$. The thermal population distribution $\Pi_{\text{th},p}$ is given by Maxwell-Boltzman distribution, and a_{in} represents the input probe field. The decay rate of photons is approximated by the free-space rate $\kappa = c/L$ with L the longitudinal extent of the atomic gas. These coupled equations describe the rich dynamics that ensue as a consequence of the collective atom-light interaction and a time-dependent pump-probe detuning.

Numerical simulations of the bistability based on this model show excellent agreement with the experimental results over a wide range of parameters of pump detuning and chirp rate (Fig. 3). In these simulations, parameters such as the pump detuning, the OD of the atomic gas, and the scan rate of the two-photon detuning were held fixed at the experimental values.

The observation of a bistable transmission requires that the momentum coherences established as a result of the atom-light interaction persist as the detuning is scanned across the RIR. In practice, off-resonant light scattering causes the thermalization of the momentum distribution thereby suppressing the bistability. This competing process leads to a limiting rate for the scan rate ($d\delta/dt$) below which the atomic momentum distribution is quasistatic and the transmission becomes nonhysteretic [Fig. 4(a)]. A bistable transmission was observed for scan rates as low as 500 kHz/ms, corresponding to moving across the RIR in 100 μs, much longer than the mean photon scattering time of <1 μs. This observation of long-lived momentum coherences is consistent with previous studies of the RIR [10,13].

An independent measure of the influence of off-resonant scattering on the probe transmission was obtained by determining the thermalization time (γ_{pop}^{-1}) of the atomic momentum distribution. For this, the probe was switched on for 100 ms at an intensity of 0.1 mW/cm^2 and at a two-photon detuning that corresponded to the peak gain of the RIR. The probe intensity was then reduced to $\sim 10^{-3}$ mW/cm^2 within a few microseconds. Following this reduction in intensity, the probe transmission indicated a gain that was initially very small but gradually relaxed to a higher equilibrium value. This relaxation was interpreted as being due to the thermalization of the momentum distribution to repopulate those states that were depleted due to the initially intense probe. Under the assumption that the final probe intensity was too small to significantly modify the atomic distribution, the time scale over which the probe transmission reaches a steady state should reflect the ther-

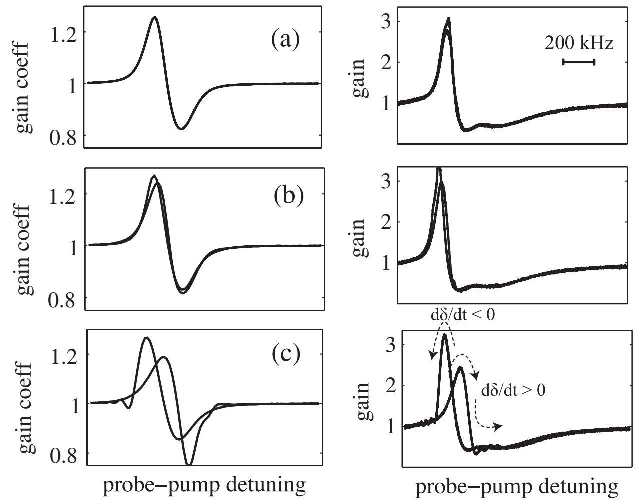

FIG. 3. Numerical simulations based on the theoretical model match the experimental results over a wide range of parameters. Panels on the left show the numerical simulations for the gain coefficient given by $\exp(-2\text{Re}[\alpha]L)$, where α is the absorption coefficient, and panels on the right show the observed experimental gain of the probe transmission across the RIR, both versus two-photon detuning. Panels (a), (b), and (c) correspond to transmission spectra obtained by chirping the two-photon detuning at scan rates of 0.1, 0.5, and 2.5 MHz/ms, respectively.

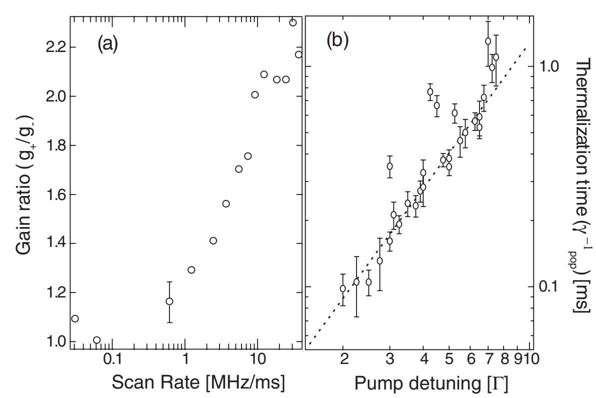

FIG. 4. (a) The ratio of peak gain (g_-/g_+) indicates a limiting scan rate $d\delta/dt \approx 0.5$ MHz/ms below which decoherence and thermalization of the momentum distribution suppress optical bistability. The data were obtained at a pump detuning $\Delta \sim -4\Gamma$. (b) Thermalization time (γ_{pop}^{-1}) vs the pump detuning. A fit to the data (dashed line) indicates a power-law dependence $\gamma_{\text{pop}}^{-1} \propto \Delta^\alpha$ with $\alpha = 1.57 \pm 0.09$.

malization time. Consistent with the observation of bistability at relatively slow scan rates, we measure thermalization times on the order of a few hundred microseconds [Fig. 4(b)].

The numerical simulations used the measured value of $\gamma_{\rm pop}$ [Fig. 4(b)] while allowing the value of $\gamma_{\rm coh}$ to vary. In order to obtain the best match with the experimental results, we typically found that $\gamma_{\rm coh}$ was required to be greater than $\gamma_{\rm pop}$. However, both quantities were much smaller than the photon scattering rate. This suppression is in agreement with earlier observations [10,13–15]. A possible explanation lies in the fact that in the presence of an optical potential ($U \propto 1/\Delta$), in the limit of large detuning, both $\gamma_{\rm coh}$ and $\gamma_{\rm pop}$ can be suppressed due to Lamb-Dicke confinement. The estimated scaling of the suppressed decay rate with the detuning should be $\gamma_{\rm eff} \propto \gamma_{\rm scatt}/\sqrt{U} \propto \Delta^{-3/2}$, where $\gamma_{\rm scatt}$ is the photon scattering rate. The experimental data [Fig. 4(b)] also suggests a similar scaling (1.57 ± 0.09).

The observation of optical bistability at low input powers suggests prospects of all-optical control using weak pulses of light. For instance, this scheme lends itself rather easily to a low-light level all-optical switch wherein the detuning of a few-photon probe is controlled in order to switch a more intense output beam. The largest scan rates, $d\delta/dt = 30$ MHz/ms, at which bistability was observed and typical values of the momentum coherence time $\gamma_{\rm coh}^{-1} \sim 100\ \mu$s together yield an estimate of the switching time $\tau = \gamma_{\rm coh}/(d\delta/dt) = 0.3\ \mu$s. This is commensurate with previous measurements of all-optical switching using a RIR [8]. Combining this with the lowest probe powers ($P \sim 20$ pW) and typical beam waists ($100\ \mu$m) used in this work, we obtain a typical photon number $\tau P/(hc/\lambda) \sim 25$ and a remarkably low switching energy density [16,17] of 7×10^{-5} photons$/(\lambda^2/2\pi)$ to operate this all-optical switch.

In conclusion, we demonstrate optical bistability due to a coherent interaction between weak light fields and the collective motion of a strongly dispersive atomic gas. Because of the relative insensitivity of the atomic momentum to external magnetic and optical fields, such strongly coupled atom-light systems are promising candidates for applications in low-light level nonlinear optics including the robust generation of correlated photons [18] and quantum memories [19]. Further, there is a close connection between the phenomenon reported here and collective atomic recoil effects such as spontaneous self-organization of the atomic medium [20], collective recoil lasing [21], and the generation of atom-photon entanglement [22].

This work was funded by the NSF, the Center for Ultracold Atoms, DARPA, and Packard foundation.

*mukundv@berkeley.edu

[1] A. Imamoglu, H. Schmidt, G. Woods, and M. Deutsch, Phys. Rev. Lett. **79**, 1467 (1997).
[2] M. D. Lukin, A. B. Matsko, M. Fleischhauer, and P. Zoller, Phys. Rev. Lett. **82**, 1847 (1999).
[3] L. M. Duan, M. D. Lukin, J. I. Cirac, and P. Zoller, Nature (London) **414**, 413 (2001).
[4] M. Fleischhauer, A. Imamoglu, and J. P. Marangos, Rev. Mod. Phys. **77**, 633 (2005).
[5] J. M. Raimond, M. Brune, and S. Haroche, Rev. Mod. Phys. **73**, 565 (2001).
[6] Th. Elsässer, B. Nagorny, and A. Hemmerich, Phys. Rev. A **69**, 033403 (2004).
[7] M. D. Lukin, Rev. Mod. Phys. **75**, 457 (2003).
[8] M. Vengalattore and M. Prentiss, Phys. Rev. A **72**, 021401 (2005).
[9] M. Vengalattore and M. Prentiss, Phys. Rev. Lett. **95**, 243601 (2005).
[10] J. Y. Courtois, G. Grynberg, B. Lounis, and P. Verkerk, Phys. Rev. Lett. **72**, 3017 (1994).
[11] J. Guo, P. R. Berman, B. Dubetsky, and G. Grynberg, Phys. Rev. A **46**, 1426 (1992).
[12] M. G. Moore and P. Meystre, Phys. Rev. A **58**, 3248 (1998).
[13] M. Kozuma, Y. Imai, K. Nakagawa, and M. Ohtsu, Phys. Rev. A **52**, R3421 (1995).
[14] P. Verkerk, B. Lounis, C. Salomon, C. Cohen-Tannoudji, J. Y. Courtois, and G. Grynberg, Phys. Rev. Lett. **68**, 3861 (1992).
[15] J. Y. Courtois and G. Grynberg, Phys. Rev. A **46**, 7060 (1992).
[16] A. M. C. Dawes, L. Illing, S. M. Clark, and D. J. Gauthier, Science **308**, 672 (2005).
[17] J. Zhang, G. Hernandez, and Y. Zhu, Opt. Lett. **32**, 1317 (2007).
[18] J. K. Thompson, J. Simon, H. Loh, and V. Vuletić, Science **313**, 74 (2006).
[19] C. Shori, B. Julsgaard, J. L. Sorensen, and E. Polzik, Phys. Rev. Lett. **89**, 057903 (2002).
[20] J. Javaloyes, M. Perrin, G. L. Lippi, and A. Politi, Phys. Rev. A **70**, 023405 (2004).
[21] D. Kruse, C. von Cube, C. Zimmermann, and Ph. W. Courtielle, Phys. Rev. Lett. **91**, 183601 (2003).
[22] N. Piovella, M. Cola, and R. Bonifacio, Phys. Rev. A **67**, 013817 (2003).