Optical bistability at low light level due to collective atomic recoil

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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 (RIR) 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.

Motivated by potential applications to quantum information science [1, 2, 3, 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], 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 EIT [6]. An added benefit of such a multiphoton process is the enhancement of the interaction time due to slow group velocities. However, since EIT 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 [7, 8], 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. Due to 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 [9]. These processes, termed ‘recoil-induced resonances’, can be described in terms of stimulated Raman transitions between different momentum classes of the atomic ensemble [10]. 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 density of the atomic gas increases, the strong atom-light coupling leads to large optical gain [7] and significant effects on the momentum distribution of the atomic gas. The strong amplification of the probe field and its subsequent back-action on the collective motional states of the gain medium results in a nonlinear optical response even for weak incident beams. In this work, this back-action 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 [7]. About 5 × 10^8 ⁸⁷Rb atoms are confined in a
Depending on the sign of the detuning chirp \((d\delta/dt)\), the transmission indicates a shift in both the resonance and the peak amplification \((g_+/g_-)\). A bistable response down to the detection limit of the input probe power is observed. The bistable response \((g_-/g_+ < 1)\) persists down to the detection limit \((\sim 20 \text{ pW})\) 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

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

where \(c_g(k)(c_e(k))\) are the annihilation operators of ground (excited) state atoms with momentum \(\hbar k\), \(\omega_0\) is the transition frequency of the two-level atoms, \(g_1(g_2)\) is the atom-light coupling coefficient and \(a_1(a_2)\) is the normalized electric field of the pump (probe) beams. In the far-detuned limit, the excited states can be adiabatically eliminated and the equation for the coherences and populations of the different momentum classes in the ground state are

\[
\frac{d}{dt} \rho(p, p') = 4i\omega_r (p'^2 - p^2) \rho(p, p') \\
+ i \frac{g_1 g_2 a_1^* \Delta}{\Delta} e^{-i\delta t} (\rho(p + 1, p') - \rho(p, p' + 1)) \\
+ i \frac{g_1 g_2 a_1 \Delta}{\Delta} e^{i\delta t} (-\rho(p, p' + 1) + \rho(p - 1, p')) \tag{2}
\]

where \(\rho(p, p') \equiv \langle c_g(k')^\dagger c_g(k) \rangle\), \(\omega_r\) 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
\[
\frac{d}{dt}a_2 = iN g_1 g_2 \frac{\Delta}{\Delta} a_1 e^{i\omega t} \sum_p \rho(p-1, p) - \frac{\kappa}{2}(a_2 - a_{in}). \tag{3}
\]

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 \(a_2\).

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

where \(\gamma_{pop}(\gamma_{coh})\) are the population (coherence) relaxation rates respectively and \(\beta = g_1 g_2 a_1 / \Delta\). The thermal population distribution \(\Pi_{th,p}\) is given by Maxwell-Boltzmann distribution. 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, OD of the atomic gas and 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 quasi-static and the transmission becomes non-hysteretic (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 \(\mu\)s, much longer than the mean photon scattering time of \(<1 \mu s\). This observation of long-lived momentum coherences is consistent with previous studies of the RIR [9, 12].

An independent measure of the influence of off-resonant scattering on the probe transmission was obtained by determining the thermalization time \(\gamma_{pop}^{-1}\) 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 thermalization 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 \(\gamma_{pop}\) (Fig. 4(b)) while allowing the value of \(\gamma_{coh}\) to vary. In order to obtain the best match with the experimental results, we typically found that \(\gamma_{coh}\) required to be greater than \(\gamma_{pop}\). However, both quantities were much smaller than the photon scattering rate. This suppression is in agreement with earlier observations [9, 12, 13, 14].
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 \( \gamma_{\text{coh}} \) and \( \gamma_{\text{pop}} \) can be suppressed due to Lamb-Dicke confinement. The estimated scaling of the suppressed decay rate with the detuning should be \( \gamma_{\text{eff}} \propto \frac{\gamma_{\text{scatt}}}{\sqrt{U}} \propto \Delta^{-3/2} \) where \( \gamma_{\text{scatt}} \) 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, \( d\delta/dt = 30 \text{ MHz/ms} \) at which bistability was observed and typical values of the momentum coherence time \( \gamma_{\text{coh}}^{-1} \sim 100 \text{ } \mu \text{s} \) together yield an estimate of the switching time \( \tau = \gamma_{\text{coh}}/(d\delta/dt) \approx 0.3 \text{ } \mu \text{s} \). This is commensurate with previous measurements of all-optical switching using a RIR [7]. Combining this with the lowest probe powers \( (P \approx 20 \text{ pW}) \) and typical beam waists (100 µm) used in this work, we obtain a typical photon number \( \tau P/(\hbar c/\lambda) \approx 25 \) and a remarkably low switching energy density \( 7 \times 10^{-5} \text{ photons}/(\lambda^2/2\pi) \) to operate this all-optical switch.

In conclusion, we demonstrate optical nonlinearities due to a coherent interaction between weak light fields and the collective motion of a strongly dispersive atomic gas. Since the atomic momentum is relatively insensitive to external magnetic or electric fields, such systems may be promising candidates for applications in low-light level nonlinear optics.

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