

Efficient and fast optical phase conjugation by use of two-photon-induced gratings in the orientation of angular momentum

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Optical phase conjugation (OPC) is of interest for many applications. The generation of squeezed light, phase-conjugate mirrors, optical correlation, and turbulence correction would all benefit from improved OPC performance. Using Zeeman sublevels and cross-circularly polarized light in rubidium vapor, we demonstrate an OPC process that uses very low power (30 mW) but is still very fast (60 ns) and efficient (gain of 20). This process is generic enough to be applicable to almost any resonant medium. © 1997 Optical Society of America

Optical phase conjugation has great potential for numerous applications. Nevertheless, current OPC materials are often limited in their usefulness. For example, photorefractives (such as BaTiO₃) have a high gain at low pump intensities (reflectivity greater than 10 at pump powers of less than 1 W/cm²). However, the response time for photorefractives (~10 ms) is not fast enough to correct for high-speed turbulence. Quantum wells can yield an extremely fast response (<1 ns) but at the cost of very low reflectivity (~0.03). Two-level atomic vapors also have fast response times (10–100 ns) with high reflectivity but only at a high pump power.¹ In this Letter we report a phase-conjugation material that has a fast response time and high reflectivity when low pump power is used, without causing overwhelming technical difficulties.

We demonstrated a high-gain optical phase conjugator in sodium vapor with a fast response time, high gain, and distortion-free propagation when low pump intensities were used.^{2,3} This double- Λ system optical phase conjugator meets most goals (reflectivity greater than 50 with a response time of less than 1 μ s at a pump power of 1 W/cm²) and was demonstrated as an effective means of turbulence correction.⁴ The double- Λ technique uses coherent population trapping (CPT) to form a grating. Recently, we observed four-wave mixing by use of CPT in a cryogenically cooled solid as well.⁵

However, for practical devices there are several problems with this double- Λ system. In the case of sodium, for example, the probe beam has to be detuned from one of the pump beams by nearly 1.77 GHz (the energy difference between the two hyperfine ground states). Furthermore, there is a strong phase coherence between these two beams such that the beat frequency is less than the natural linewidth of the sodium D1 transition. Moreover, only bulky dye lasers

can produce the 589-nm wavelength at the necessary power. Other alkali vapors (such as Rb and Cs) have similar transitions that one can access by diode lasers, but the problem of a large detuning (6.8 GHz for ⁸⁷Rb, 9.2 GHz for Cs) with sharp phase coherence remains. For practical systems one would prefer to have more flexibility in choosing the optical wavelengths.

Here we produce Λ systems that use Zeeman sublevels in a vapor cell, which allows a single hyperfine transition to contain the Λ systems. In the four-wave-mixing process forward pump F and signal beam S produce a Zeeman sublevel coherence grating by CPT. Backward pump B diffracts off this grating, forming the conjugate beam. The polarizations are then critical for ensuring the formation of the Λ systems. One obvious advantage of this approach is that there is no large detuning between the pump beams. Another significant advantage is that almost any transition in almost any resonant system can work. Since typical hyperfine ground-state splitting is large compared with the Maxwell–Boltzmann Doppler width (of a single ground-state transition) in a vapor cell, atoms in the unused hyperfine ground state are transparent. The natural repumping processes in a cell (e.g., collisions) can provide a continuous sample of atoms in the active hyperfine ground state. We should point out that coherences between Zeeman split ground-state levels were used in a variety of experiments^{6–8} but not for the process reported here.

To help provide a physical picture of the coherence grating, we consider a $J = 1$ to $J' = 1$ transition with no magnetic field present. We now apply a forward beam and a signal beam with degenerate frequencies and equal Rabi rates but opposite circular polarizations. We also assume that the forward and the signal beams have a small angle difference, 2α . Since the $J' = 1$, $m_z' = 0$ state does not decay into the $J = 1$, $m_z = 0$

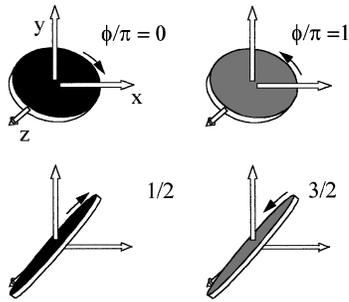


Fig. 1. In the properly rotated frame there is an $m = 0$ state with all of the atomic population. The orientation of this $m = 0$ state varies with the grating phase, ϕ (see text). The axes shown are in the laboratory frame coordinates.

transition (a forbidden dipole transition), all the atoms are rapidly optically pumped into a dark state that is a superposition of the $m_z = 1$ and $m_z = -1$ states (owing to CPT). Specifically, the dark state is given by

$$|D\rangle = \frac{1}{\sqrt{2}}[|m_z = 1\rangle + \exp(i\phi)|m_z = -1\rangle],$$

$$\phi = \phi(x) \equiv (\mathbf{k}_F - \mathbf{k}_S) \cdot \mathbf{r}$$

We define our laboratory frame-coordinate systems such that

$$\mathbf{k}_F = k[\sin(\alpha)\hat{x} + \cos(\alpha)\hat{z}],$$

$$\mathbf{k}_S = \mathbf{k}[-\sin(\alpha)\hat{x} + \cos(\alpha)\hat{z}].$$

For a given value of x it is possible to find a quantization axis \tilde{z} such that the dark state becomes $|D\rangle = |m_{\tilde{z}} = 0\rangle$. To help illustrate the grating we use the vector model for azimuthal spin. This pictorial model uses a vector of length $[j(j+1)]^{1/2}\hbar$, which precesses about the \tilde{z} axis. The projection of this vector onto the \tilde{z} axis determines $m_{\tilde{z}}$. We note that the $m_{\tilde{z}} = 0$ state is disk shaped, with \tilde{z} normal to the disk. Figure 1 shows how the orientation of this dark-state disk varies with $\phi(x)$, thus forming a grating in the x direction.

In this Letter we concentrate on the $^{87}\text{Rb } 5^2S_{1/2} F = 2 \rightarrow 5^2P_{1/2} F = 1$ transition (see Fig. 2). This transition yielded the highest phase-conjugate reflectivity; however, we also observed robust conjugates for other transitions. The frequency and the polarization of each beam are indicated, and for simplicity only one Λ system is shown.

Figure 3 demonstrates how we produce the beams for this experiment. All beams originate from a single Ti:sapphire laser. We produce F , B , and S by use of acousto-optic modulators and beam splitters. Wave plates and polarizing beam-splitting cubes are used to get the appropriate polarizations. The four-wave mixing occurs in a heat-pipe oven (Rb cell) at an approximate temperature of 140°C . Helmholtz coils produce the magnetic field and mu-metal encases both the Helmholtz coils and the Rb cell. Typically the two pump beams have a 1-mm FWHM, whereas the signal has a 0.6-mm FWHM. All beams cross with angles below 5 mrad.

Figure 4 demonstrates the frequency dependence of the conjugate as we scan the Ti:sapphire laser from the $^{87}\text{Rb } 5^2S_{1/2} F = 2 \rightarrow 5^2P_{1/2} F = 1$ transition to the $5^2P_{1/2} F = 2$ transition (these transitions are marked 1

and 2, respectively, in Fig. 4). F has a power of 14 mW, B has a power of 8 mW, $\Delta = 81$ MHz and $\mathbf{B} = 1.50$ G. We found that the conjugate is almost entirely σ_+ polarized as expected (the σ_+ component is ~ 400 times stronger than the σ_- component at the peak). At the frequency corresponding to the highest conjugate power, the reflectivity is 3. When we increase the pump power by a factor of 4, we get a reflectivity of close to 300 (not shown).

Figure 5 shows the response of conjugate C for a changing magnetic field (e.g., Δ_z). Since the Zeeman

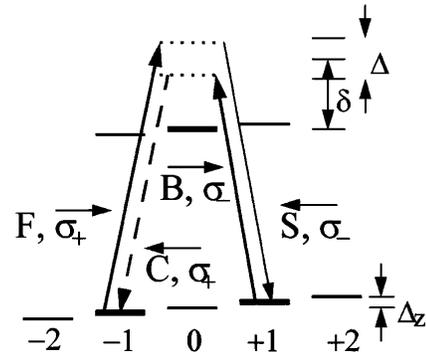


Fig. 2. Four-wave-mixing interaction in the $^{87}\text{Rb } 5^2S_{1/2} F = 2 \rightarrow 5^2P_{1/2} F = 1$ transition. Only one of the possible Λ systems is shown.

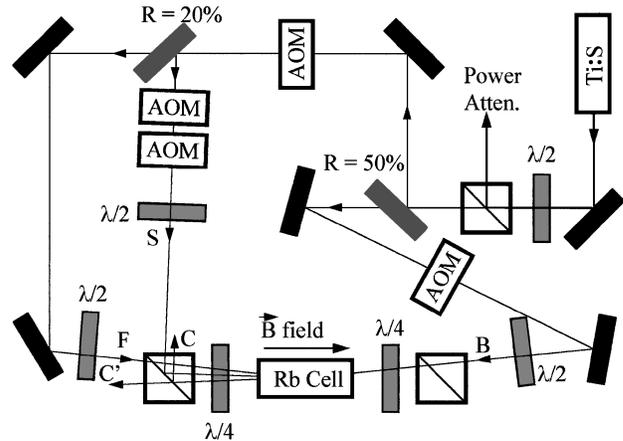


Fig. 3. Schematic illustration of the experimental setup for producing phase conjugates with a heat-pipe oven. AOM's, acousto-optic modulators. Atten., attenuation.

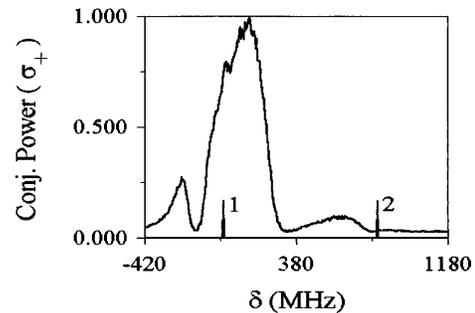


Fig. 4. Dependence of the phase-conjugation process on the laser frequency. The conjugate (Conj.) power is in arbitrary units.

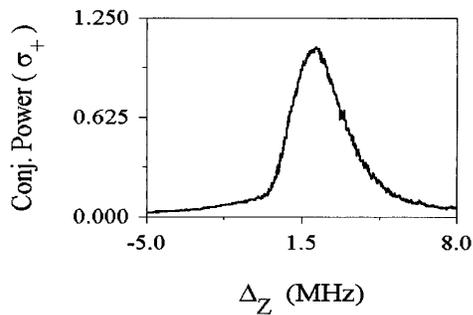


Fig. 5. Dependence of the conjugate power on the magnetic field (equivalent, in this case, to two-photon detuning).

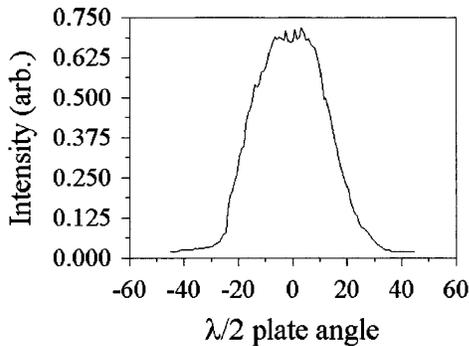


Fig. 6. Dependence of the conjugate power on the polarization of the backward pump.

shift for small fields is well understood, we can find the FWHM for this scan (2.4 MHz). This subnatural linewidth ($\Gamma = 5.89$ MHz) is quite consistent with our proposed mechanism. By examining the system (see Fig. 2), one can see that a change in the magnetic field is similar to a change in the signal detuning. As a self-consistency check we also detuned the signal beam (not shown). By scanning this way, we find a FWHM of 3.4 MHz, which is slightly larger than the linewidth ow-

ing to direct two-photon detuning. We attribute this increase in width to spatial inhomogeneity in the magnetic field. If we assume that the linewidth primarily is due to the optical pumping rate, we find a response time of ~ 60 ns.

Finally, we examined the conjugate's dependence on the backward pump polarization as shown in Fig. 6. By inserting a $\lambda/2$ plate between the beam-splitting cube and the $\lambda/4$ plate, we can cycle through all polarizations. When the $\lambda/2$ plate is rotated 45° in either direction, the backward beam is changed from σ_- to σ_+ polarization. Thus, as our model would predict, the conjugate beam disappears when the polarization is incorrect.

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