Optics Communications 95 (1993) 255–259
North-Holland

OPTICS
COMMUNICATIONS

Polarization oscillations of coupled laser beams in an optically pumped atomic vapour

Dieter Suter
Institute of Quantum Electronics, Swiss Federal Institute of Technology (ETH) Zürich, CH-8093 Zürich, Switzerland

Received 7 August 1992

Spontaneous polarization oscillations are observed to develop in two laser beams interacting in an atomic vapour when their polarizations are locked to each other. In a theoretical analysis of the situation, we show that these oscillations are due to ground-state magnetization undergoing Larmor precession. The gain of the feedback loop that locks the two polarizations acts as a control parameter that can cause a phase-transition of the combined system (atoms and laser beams) into a stable limit cycle.

1. Introduction

The large optical nonlinearities that are produced by optical pumping in atomic vapours have often been used to study nonlinear dynamics. As an example, it was shown that counterpropagating laser beams in atomic sodium, tuned close to the transition frequency of the D_1 resonance line, can show amplitude- [1] and polarization- [2] instabilities. Self-pulsing and temporal chaos in the polarization of the transmitted light were also observed when the atomic medium was placed in an optical cavity and a magnetic field was applied perpendicular to the direction of the laser beam [3,4]. In these experiments, it was shown that the period of the self-pulsing period can be controlled by the strength of a magnetic field.

On a qualitative level, the self-pulsing can be understood as originating from a precessing magnetic moment which is induced in the electronic ground state of the atomic medium by the optical pumping process and forced into Larmor precession by the magnetic field. This polarization of the angular momentum substates leads to a long-lived optical anisotropy which modulates the polarization of a transmitted laser beam [5,6]. On the other hand, precessing ground state magnetization can be produced very efficiently by modulated optical pumping [7–9]. The feedback of transmitted light into the atomic medium by the optical cavity leads therefore to a nonlinear coupling between atoms and radiation and, for a range of control parameters, to instabilities like self-pulsing or chaos.

It this Communication, we want to demonstrate an experimental setup that allows the observation of instabilities like self-pulsing at considerably lower laser intensities. This is achieved by using two separate laser beams whose polarizations are locked to each other. One of the two beams acts as an optical pump beam while the other serves as a probe; the polarization of the latter is analysed behind the sample in a polarization-selective detector and the resulting signal is used to control the polarization of the pump beam via an electronic feedback loop. This polarization-feedback allows one to cancel the damping of the ground state magnetization without affecting the Larmor precession. The resulting behaviour of the atoms and laser beam polarizations depends on the strength of the coupling between the two beams: for weak coupling, i.e. for low loop gains, the free precession time is simply extended, but for strong coupling, the system falls into a stable limit cycle with freely precessing magnetization and oscillating polarization of the coupled laser beams.

2. Theory

The proposed setup is shown schematically in fig. 1: an optical pump beam (full line) excites population differences in an atomic medium (sodium in the experimental example). A second laser beam (dashed line) probes this population difference; by passing it through the polarization-selective detector shown in the figure, a signal is obtained that is directly proportional to the population difference [10]. Via an amplifier, this signal is fed into an electrooptic modulator (EOM) whose fast axis is oriented at 45° with respect to the (linear) polarization of the incident pump beam. After passing through the EOM, the pump beam polarization is in general elliptical, with the ellipticity depending on the modulator voltage; it becomes circular when the amplified signal equals the quarter-wave voltage of the EOM.

For the description of the resulting dynamics, we use a coordinate system whose z-axis is parallel to the laser beams and whose x-axis is given by the magnetic field direction. We assume that the atomic medium can be modelled as a $J=1/2 \leftrightarrow J'=1/2$ transition and that the laser intensity is low enough that excited state effects can be neglected. With these assumptions, we can describe the ground state dynamics in terms of the three-component magnetization vector $\boldsymbol{m}=(m_x, m_y, m_z)$ whose z-component represents the ground state population difference while the x- and y-components correspond to coherences between the Zeeman sublevels [3]. The optical pump rates $P_\pm$ for the population pumping into the $m_J=\pm 1/2$ ground state sublevels are then proportional to the intensities of the circular polarization components. Since the overall intensity does not change during the modulation cycle, we have [11]

$$P_+ + P_- = P_0 , \quad P_+ - P_- = P_0 \sin(g m_z) , \tag{1}$$

where P_0 is the pump rate of a circularly polarized beam. The dimensionless gain factor g, which describes the coupling between the two laser beam polarizations, includes the sensitivity of the detector, which produces a signal proportional to m_z, the gain of the electronic feedback loop and the electrooptic coefficient of the modulator. Neglecting light-shift effects, the equation of motion for the atomic system can be written as

$$\begin{aligned} \dot m_y &= +\Omega_L m_z - (\gamma_0 + P_+ + P_-) m_y , \\ \dot m_z &= P_+ - P_- - \Omega_L m_y - (\gamma_0 + P_+ + P_-) m_z , \end{aligned} \tag{2}$$

where γ_0 describes the loss of magnetization due to diffusion. The x-component is not coupled to the other two variables and is not considered here. Eliminating the optical pump rates $P_\pm$, we can combine eqs. (1), (2) into an equation of motion for the atomic medium alone,

$$\begin{aligned} \dot m_y &= +\Omega_L m_z - (\gamma_0 + P_0) m_y , \\ \dot m_z &= -\Omega_L m_y - (\gamma_0 + P_0) m_z + P_0 \sin(g m_z) . \end{aligned} \tag{3}$$

As these equations do not have an analytic solution, we discuss first the regime $g m_z \ll 1$, where they reduce to

$$\begin{aligned} \dot m_y &= \Omega_L m_z - (\gamma_0 + P_0) m_y , \\ \dot m_z &= -\Omega_L m_y + [P_0(g-1) - \gamma_0] m_z . \end{aligned} \tag{4}$$

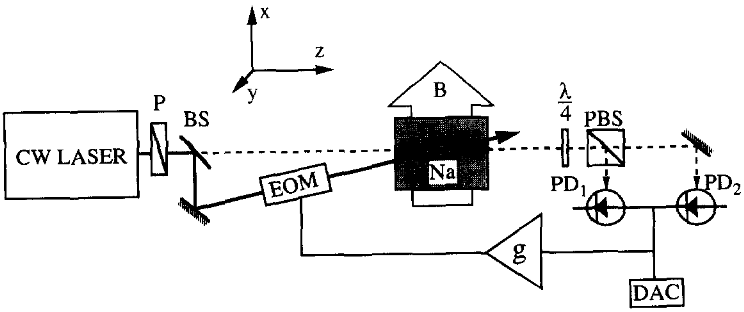

Fig. 1. Schematic representation of the experimental setup. P = polarizer, BS = beam splitter, EOM = electrooptic modulator, B = magnetic field, $\lambda/4$ = retardation plate, PBS = polarizing beam-splitter, $PD_{1,2}$ = photodiodes, DAC = D/A converter.

We are specially interested in the solutions for the case of small gain, $g \ll \Omega_L/P_0$:

$$m_y(t) = m_0 \sin(\Omega_L t - \phi) \exp(kt) ,$$
$$m_z(t) = m_0 \cos(\Omega_L t - \phi) \exp(kt) , \quad (5)$$

where the starting amplitude m_0, and the phase ϕ are determined by the initial conditions. The sign of the growth rate $k = [P_0(g/2-1) - \gamma_0]$ determines whether the magnetization in the system decays to zero or grows exponentially. The case of small or vanishing gain $(0 \leqslant g < 2)$ is equivalent to freely precessing magnetization, with a damping rate that is increased by the term $P_0(1-g/2)$ due to the damping effect of the partially polarized pump beam. For $g=2$, this additional damping effect by the pump beam is cancelled exactly, and in the region $2 < g < 2(1+\gamma_0/P_0)$, the FID decays with a time constant that is longer than during free precession and diverges for $g = 2(1+\gamma_0/P_0)$.

If the gain is increased beyond this critical value, the magnetization does no longer decay, and the system falls into a stable limit cycle where the magnetization processes at the Larmor frequency with a constant amplitude. The value of this amplitude cannot be calculated from the linearized equations (4). We consider therefore an equation of motion for the amplitude $m_{yz} = \sqrt{m_y^2 + m_z^2}$ alone which is valid as long as $k \ll \Omega_L$, i.e. when the amplitude changes only slightly during a Larmor cycle:

$$\dot{m}_{yz} = -(\gamma_0 + P_0) m_{yz} + P_0 J_1(g m_{yz}) , \quad (6)$$

where J_1 represents the first order Bessel function. The two terms describe the competing effects discussed above.

The magnetization amplitude of the limit cycle is determined by the stationary solutions of eq. (6) which are displayed as a function of g in fig. 2 for the parameter values $\gamma_0/P_0 = 0.1$. The highest amplitude is apparently reached for a gain factor $g \sim 3.5$; increasing the gain further lowers the amplitude because the increase in the voltage across the EOM raises the excitation efficiency at higher harmonics of the Larmor frequency at the cost of the fundamental which is the only one contributing to the buildup of magnetization [11]. Multiple solutions are possible if the gain is increased further; these additional solutions reflect the oscillatory nature of the Bessel function. The solutions corresponding to stable limit cycles are drawn as bold lines in fig. 2, while those corresponding to unstable trajectories are drawn as thin lines.

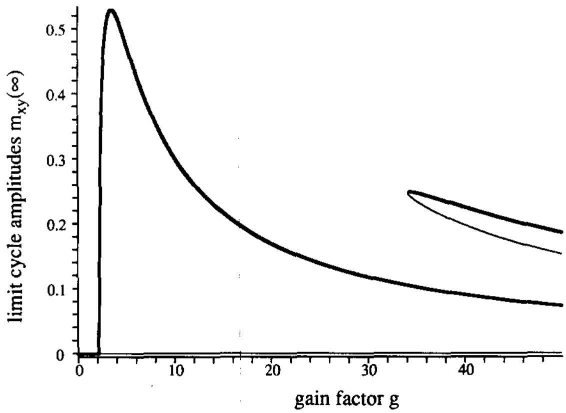

Fig. 2. Magnetization amplitude during the limit cycles as a function of the gain in the feedback loop for $P_0/\gamma_0 = 10$. The thick lines denote stable limit cycles, the thin lines unstable ones.

3. Experimental results and discussion

The setup used for the experimental implementation of this scheme is shown schematically in fig. 1. The atomic medium, sodium in our case, was contained in a glass cell at a temperature of 130 °C. In order to have a homogeneous optical resonance line, 130 mbar of argon was added as a buffer gas; the collisional broadening caused a homogeneous linewidth of 1.5 GHz, somewhat larger than the Doppler width. The laser beams were derived from a single-mode cw ring dye laser; the laser was tuned to the Na-D_1 line ($\lambda = 589.6$ nm) where the small signal absorption was $\sim 60\%$. The output of the laser was split into a pump beam with an intensity of 200 μW/mm^2 and a probe beam with an intensity of 50 μW/mm^2 which overlapped in the sample region at an angle of less than 1°. The polarization of both incident laser beams was linear and the orientation of the pump laser polarization was set to 45° with respect to the fast axis of the EOM. The static birefringence of the modulator crystal was compensated with a Soleil–Babinet compensator. Behind the sample cell, the pump beam was blocked and the probe beam was passed into a polarization-selective detector [10] which measured the intensity difference of the two circular polarization components, so that the resulting signal was proportional to the ground-state magnetization component m_z. This signal was amplified and fed into the EOM, thus coupling the polarization of the pump beam to that of the probe beam. At the same time, the detector signal was digitised and transferred to a computer for data analysis and storage.

The experiments presented here were performed in a magnetic field of $|B_x| = 14.3$ μT, corresponding to a Larmor frequency of 100 kHz. The feedback loop was closed by switching the pump laser beam on with an acoustooptic modulator (not shown in fig. 1). As displayed in fig. 3, the polarization of the laser beams starts to oscillate soon after the pump laser beam is switched on; the oscillations grow out of the noise with a delay that depends on the loop gain. The phase varies randomly between experiments, as expected for a random initial condition. The growth of the oscillation amplitude is initially roughly exponential until the system saturates and reaches the limit cycle. When the pump laser is switched off, the magnetization, which was built up during the pulse, is observed as a free induction decay. Comparing the oscillation period of the FID with that of the driven magnetization shows that the driven system does indeed oscillate at the Larmor frequency. A closer examination of the equations of motion shows that

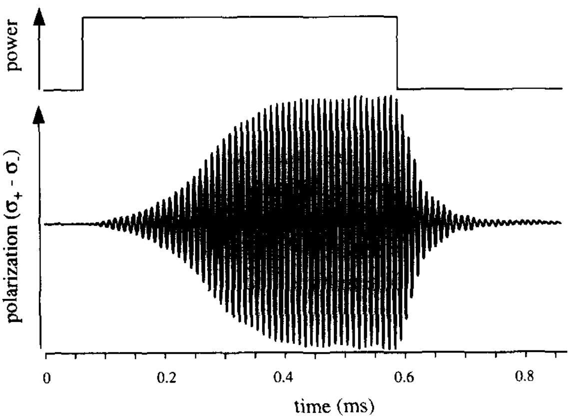

Fig. 3. Polarization oscillation observed when the feedback loop is closed by switching the pump laser on. The upper trace shows the pump laser intensity, the lower trace the signal of the polarization-selective detector.

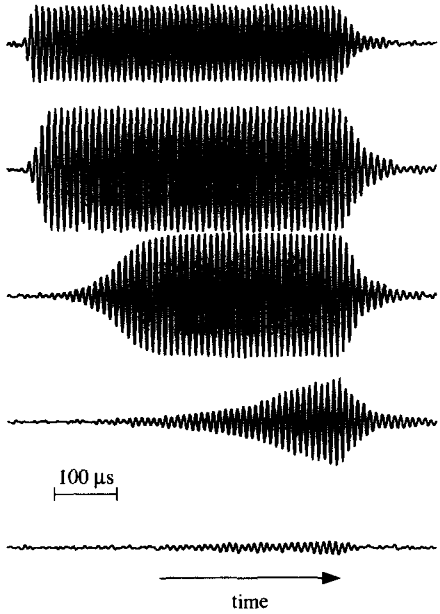

Fig. 4. Polarization oscillations as a function of the loop gain which increases from bottom to top.

propagation delays through the feedback loop must be negligible if the oscillation frequency is to be determined by the magnetic field strength alone.

Figure 4 shows a series of self-starting oscillations for increasing gain (from bottom to top) in the electronic feedback loop. In the lowest trace, the gain is just high enough to cancel the losses and the growth of the oscillations is very slow. For increasing gain, the amplitude as well as the rate of the approach to the stationary state increase until the amplitude reaches a maximum. For the highest gain (top trace), the rate increases further, but the amplitude decreases again, in agreement with the theoretical curves shown in fig. 2.

4. Conclusion

We have demonstrated an experimental setup that makes it possible to lock the polarizations of two laser beams to each other that interact with the same atomic medium. We have analysed the resulting dynamics of the atomic medium and the coupled laser beam polarizations; it was found that the system can, under suitable conditions, undergo a spontaneous symmetry breaking and develop a magnetization which undergoes Larmor precession if a transverse magnetic field is applied. Depending on the experimental parameters, such as the loop gain, the precessing magnetization can be strongly damped, undamped or grow exponentially until it falls into a stable limit cycle which corresponds to a continuously precessing magnetization vector of constant length and synchronously oscillating polarizations of the laser beams. Numerical simulations, using the equations of motion presented here, have shown that the system can also exhibit period doubling and chaotic motion. Since the experimental setup is simple and rather robust, this may become an interesting prototype system for studying nonlinear dynamics.

The results reported here can be compared to the 'NMR-laser' of Brun and coworkers [12,13], where a polarized system of nuclear spins is placed in a coil that is a part of a tuned rf-circuit acting as a laser cavity. Like in the present case, the dynamics of the NMR laser are determined by the coupling between the spin system and the magnetic field. In our setup, the locking of the polarizations of the two beams may be considered as one mirror, while the interaction of the beams in the atomic medium represents the second mirror of the laser cavity.

Since the system locks itself to the Larmor frequency of the ground state magnetization, it can be considered as a magnetic field controlled oscillator, which converts the field strength into a frequency. As such, it could be used as a sensitive magnetometer. Compared to other setups [14], it should provide similar absolute sensitivity, but the relative sensitivity could be significantly higher, since the value of the bias field does not influence the absolute sensitivity, as long as the delays in the feedback circuit are negligible on the timescale of one Larmor period.

Acknowledgements

This work was supported by the Schweizerischer Nationalfonds.

References

[1] G. Khitrova, J.F. Valley and H.M. Gibbs, Phys. Rev. Lett. 60 (1988) 1126.
[2] D.J. Gauthier, M.S. Malcuit and R.W. Boyd, Phys. Rev. Lett. 61 (1988) 1827.
[3] F. Mitschke, R. Deserno, W. Lange and J. Mlynek, Phys. Rev. A 33 (1986) 3219.
[4] C. Boden, M. Dämmig and F. Mitschke, Phys. Rev. A 45 (1992) 6829.
[5] H.G. Dehmelt, Phys. Rev. 105 (1957) 1924.
[6] S. Pancharatnam, Phys. Lett. A 27 (1968) 509.
[7] W.E. Bell and A.L. Bloom, Phys. Rev. Lett. 6 (1961) 280.
[8] J. Mlynek, K.H. Drake, G. Kersten, D. Fröhlich and W. Lange, Optics Lett. 6 (1981) 87.
[9] D. Suter and J. Mlynek, Phys. Rev. A 43 (1991) 6124.
[10] D. Suter, M. Rosatzin and J. Mlynek, Phys. Rev. A 41 (1990) 1643.
[11] H. Klepel and D. Suter, Optics Comm. 90 (1992) 46.
[12] H. Marxer, B. Derighetti and E. Brun, Phys. Rev. Lett. 50 (1983) 958.
[13] E. Brun, B. Derighetti, D. Meier, R. Holzner and M. Ravani, J. Opt. Soc. Am. B 2 (1985) 156.
[14] C. Cohen-Tannoudji, J. Dupont-Roc, S. Haroche and F. Laloë, Rev. Phys. Appl. 5 (1970) 102.