Reversible self–Kerr nonlinearity in an N-type atomic system through a switching field

Xu Yang1, Kang Ying2, Yueping Niu1 and Shangqing Gong1

1 Department of Physics, East China University of Science and Technology, Shanghai 200237, People’s Republic of China
2 Shanghai Institute of Optics and Fine Mechanics, Chinese Academy of Sciences, Shanghai 201800, People’s Republic of China

E-mail: niuyp@ecust.edu.cn

Received 28 November 2014, revised 2 February 2015
Accepted for publication 24 February 2015
Published 27 March 2015

Abstract
We investigate the self–Kerr nonlinearity of a four-level N-type atomic system in 87Rb and observe its reversible property with a unidirectional increase in the switching field. For a laser arrangement in which the probe field interacts with the middle two states, the slope and the sign of the self–Kerr nonlinearity around the atomic resonance not only can be changed from negative to positive but also can be changed to negative again with a unidirectional increase in the switching field. Numerical simulations are consistent with the experimental results, and dressed state analysis is presented to explain the experimental results.

Keywords: quantum optics, Kerr effect, nonlinear

1. Introduction
Kerr nonlinearity, which can be used in quantum non-demolition measurements [1, 2], quantum logic gates [3, 4], single-photon nonlinearity [5], and the generation of optical solitons [6, 7], has been a research hotspot in recent years. Many approaches have been taken to achieve significant enhancement of Kerr nonlinearity [8–13], among which the use of electromagnetically induced transparency (EIT) [14, 15] is an important technology because of reduced absorption. In addition to enhancement, sign and slope modification of Kerr nonlinearity also have specific applications in many fields, such as all-optical switches [16–18] and logic gates [19, 20] in optical communications and quantum computers. Recently the Xiao group changed the slope and sign of the self–Kerr nonlinearity from negative to positive just by increasing the power of the additional switching laser [21]. In this work, we report our investigation of the self–Kerr nonlinearity in a similar N-type four-level system but with a different arrangement of laser fields. With a unidirectional increase in the switching laser power, the slope and sign of the Kerr nonlinearity not only can be changed from negative to positive but also can be changed to negative again. Such a change of the sign or slope of the nonlinear coefficient can be used to form a more complicated optical system. Dressed state analysis is presented to demonstrate the foregoing experimental observations.

2. Experimental arrangement
The N-type system used in our experiment is the D2 line (780 nm) of 87Rb. The hyperfine states of 5S1/2 (F = 1 and F = 2) and 5P3/2 (F′ = 1 and F′ = 3) are used to form a four-level N-type configuration. Unlike the laser field arrangement in reference [21], here the coupling field Ωc interacts with the ground state |1⟩ and the excited state |3⟩, whereas the probe field Ωp is applied between states |2⟩ and |3⟩. One may note that here spontaneous emission from |4⟩ to |1⟩ is dipole forbidden, so it will not show gain, in accordance with reference [22].

Figure 2 is our experimental setup, which is similar to the one used in our previous experiments [23–25]. The coupling, probe, and switching fields are all single-mode tunable external cavity diode lasers (ECDLs) (New Focus TLB-6900) with a linewidth of about 300 kHz. AOM is the acousto-optic modulator, which can adjust the frequency of the probe field. The half-wave plate 1 (HWP1) and polarized beam splitter 1
probe detuning, and switching fields. Here $\Delta_p = \omega_{p2} - \omega_p$ is the probe detuning, and $\Delta_c = \omega_{c1} - \omega_c$, $\Delta_s = \omega_{s2} - \omega_s$ are the coupling and switching field detunings.

Before entering the cavity, the coupling and switching beams are brought together by PBS2 to interact with the Rb atoms. The coupling and switching laser power adjusted by HWP2, HWP3, and HWP4 is separated into an auxiliary Rb cell for stabilizing the frequency of the two beams by using the saturated absorption spectroscopy (SAS) method. Then the three beams are brought together by PBS2 to interact with the Rb atoms. Before entering the cavity, the coupling and switching beams are focused by lenses with a focal length of 30 cm to make their beam diameters approximately 300 $\mu$m. PBS5 is used to reject the switching and coupling fields. The reflectivity of the flat mirror M1 is approximately 99.5%. The cavity mirror M2, with a reflectivity of 99.5%, is concave with a 15-cm radius of curvature and is controlled by a piezoelectric (PZT) driver. The finesse of the empty cavity (without the Rb cell and PBS2) is about 138. After we insert the Rb cell and PBS2, the finesse of the cavity (with the Rb atoms off resonance) is reduced to 66 because of surface reflection losses.

In our experiment, first we use the SAS method to lock the frequency of the coupling and switching fields to be resonant with the corresponding transitions. As for the probe field, its frequency then can be tuned and locked by adjusting the AOM and SAS. Cavity scanning is carried out by tuning the applied voltage of the PZT that is attached to the cavity mirror M2, and the scanning rate is about 5.47 MHz ms$^{-1}$.

When the cavity is scanned, the transmission spectrum shows a typical symmetric Lorentzian shape because the probe beam is tuned far from the atomic resonance. But when the probe field is tuned near the atomic resonance, the cavity transmission profile becomes asymmetric due to the self–Kerr nonlinearity. Figure 3 shows the cavity transmission when $\Delta_p = 2$ MHz, $\Delta_c = \Delta_s = 0$, the coupling laser power $P_c = 0.9$ mW, and the switching laser power $P_s = 0.7$ mW. According to the relation between the nonlinear phase shift of the optical cavity and the nonlinearity of the inside medium, the self–Kerr nonlinear coefficient $n_2$ can be directly obtained from the asymmetry degree of the cavity transmission [26, 27]. Hence, we can measure the coefficient $n_2$ by scanning the cavity length through the PZT for each determined probe frequency. The probe field is totally tuned in the range of $\pm$50 MHz to 50 MHz. For a different probe frequency, the intracavity power will change due to different absorption, so we should adjust the input probe laser power sightly in our experiment.

3. Comparison and discussions

The left column in figure 4 shows our experimentally measured self–Kerr nonlinear coefficient $n_2$ at different probe detunings. From (a1) to (a5), $P_s$ is increased from 0.3 mW to 2 mW, whereas $P_c$ is maintained at 0.9 mW. When $P_s$ is much smaller than $P_c$, as shown in (a1) and (a2), the slope of $n_2$ near resonance is negative. But as $P_s$ reaches 0.9 mW, the slope of $n_2$ near resonance changes from negative to positive (figure 4(a3)). This phenomenon is similar to that in reference [21]. But what is worth pointing out is that as $P_s$ increases further, the slope of $n_2$ near resonance goes back to being negative, which is different from the phenomenon in reference [21]. As displayed in figures 4(a4) and (a5), for the cases with 1.2 mW and 1.5 mW, the slope of $n_2$ near resonance is negative.

Based on the density-matrix equations that describe the motion of the atoms under the interaction of the three laser fields, we get the expression of self–Kerr nonlinearity through the steady-state solution under the weak-probe approximation [21, 28]. Also, we have considered the Doppler effect and collision decay in our numerical simulation. The right column in figure 4 represents our theoretical numerical result. One can...
see that the slope of $n_2$ near resonance also changes from negative to positive and again to negative with the unidirectional increase in the switching field power. This trend is consistent with our experimental result. But one may note that the magnitude of $n_2$ in (b3) is much larger than the other calculation results, whereas the corresponding experimental magnitude is not so high. This inconsistency is caused by at least the following two factors. First, from (b3) we can see

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig4}
\caption{Left column: the experimental measured self–Kerr nonlinear coefficient $n_2$. $P_c = 0.9$ mW. From (a1) to (a5), $P_s = 0.3$ mW, 0.7 mW, 0.9 mW, 1.2 mW, 1.5 mW. Right column: the calculated self–Kerr nonlinear coefficient $n_2$ with the parameters $\Omega_\pi = 2\pi \times 35$ MHz and $\Omega_s = 2\pi \times 12$ MHz, $\Omega_\eta = 2\pi \times 24$ MHz, $\Omega_\varsigma = 2\pi \times 35$ MHz, $\Omega_\xi = 2\pi \times 47$ MHz, $\Omega_\zeta = 2\pi \times 58$ MHz, $\gamma_1 = \gamma_2 = \gamma_3 = 2\pi \times 3$ MHz, $\gamma_4 = 0$, where $\gamma$ is the decay rate of the energy levels in figure 1. The collision decay is $\gamma^{col} = 2\pi \times 6$ MHz. $T = 28^\circ C$.}
\end{figure}
that the curve of \( n_2 \) has a very sharp change in a very small range of \( \Delta_p \). However, the measurement step we used in the experiment is 3 MHz, so the position we measured is outside the peak value. Second, according to our simulations (as one can see figure 5), the value of \( n_2 \) has a sharp decrease when \( \Omega_c \) is not equal to \( \Omega_e \). In contrast, experimentally it is hard to keep \( \Omega_c = \Omega_e \) because the effective interaction \( (\Omega_c/\Omega_e) \) is determined by several parameters, including light spot size and dipole moment. Therefore, the deviations of these factors may cause \( \Omega_e \) to be not exactly equal to \( \Omega_c \) in our experiment when \( P_s = P_c \). As a result, there is a difference between (a3) and (b3) in figure 4.

When it comes to the sign of \( n_2 \), we take the value at about \( \Delta_p = 2 \text{ MHz} \) as an example. Just as figure 5 shows, the red spot is the experimental value of the self–Kerr nonlinear coefficient \( n_2 \) and the black spot is the corresponding theoretical value in the same position of resonance. We can see that \( n_2 \) changes from negative to positive and then reverses to negative with the unidirectional increase in the switching laser power. The experimental and theoretical results are basically consistent.

4. Dressed state analysis

In this section, we present the dressed state analysis of this N-type system. For a weak probe field, levels \( |1\rangle \) and \( |3\rangle \) are dressed by \( \Omega_e \) and levels \( |2\rangle \) and \( |4\rangle \) are dressed by \( \Omega_c \) separately. The corresponding dressed states can be represented with

\[
|\alpha_{c1}\rangle = \sqrt{2} \left( |3\rangle + |1\rangle \right) / 2 \\
|\alpha_{c2}\rangle = \sqrt{2} \left( |3\rangle - |1\rangle \right) / 2 \\
|\alpha_{s1}\rangle = \sqrt{2} \left( |4\rangle + |2\rangle \right) / 2 \\
|\alpha_{s2}\rangle = \sqrt{2} \left( |4\rangle - |2\rangle \right) / 2
\]

The energy difference between the eigenvectors \(|\alpha_{c1}\rangle \rightarrow |\alpha_{c2}\rangle\) and \(|\alpha_{s1}\rangle \rightarrow |\alpha_{s2}\rangle\) are \( \Omega_e \) and \( \Omega_c \) respectively, just as figure 6(a) shows. Apparently, regardless of whether \( \Omega_c \) is larger or smaller than \( \Omega_e \), there are always four absorption peaks. Thus, transparency occurs at the probe’s resonant position, where the slope of the Kerr nonlinear coefficient \( n_2 \) is negative. However, when \( \Omega_c \) is equal to \( \Omega_e \), we can see only three absorption peaks. That is to say, the middle two absorption peaks are simply combined into one, so the system becomes absorptive at the probe’s resonant position, where the slope of \( n_2 \) is positive. These results correspond to the observed experimental phenomenon shown in figure 4(a). Considering the natural linewidth and the applied field, the absorption peak will have a certain width. So when \( P_s \) is close to \( P_c \), the slope of \( n_2 \) near resonance is positive, and for other \( P_s \) and \( P_c \), it is negative.

If we use the foregoing dressed state analysis to inspect the situation of the other laser arrangement in reference [21], we find that when \( \Omega_e \) is small, it only slightly disturbs the EIT and the slope of \( n_2 \) near resonance is still negative. As \( \Omega_c \) increases, it dresses the corresponding energy levels, along with \( \Omega_e \). Then the three dressed states can be written as [29]

\[
|\beta_0\rangle = \left( -\Omega_s |3\rangle + \Omega_c |4\rangle \right) / \sqrt{\Omega_c^2 + \Omega_s^2} \\
|\beta_1\rangle = \left( |2\rangle + \left( \Omega_c / \sqrt{\Omega_c^2 + \Omega_s^2} \right) |3\rangle \right) + \left( \Omega_c / \sqrt{\Omega_c^2 + \Omega_s^2} \right) |4\rangle \bigg) / \sqrt{2} \\
|\beta_2\rangle = \left( -|2\rangle + \left( \Omega_c / \sqrt{\Omega_c^2 + \Omega_s^2} \right) |3\rangle \right) + \left( \Omega_c / \sqrt{\Omega_c^2 + \Omega_s^2} \right) |4\rangle \bigg) / \sqrt{2}
\]

The energy difference between the three eigenvectors is \( \sqrt{\Omega_c^2 + \Omega_s^2} \) in both cases, as figure 6(b) shows. Hence the system is absorptive at the probe’s resonant position. Therefore, the slope of \( n_2 \) simply changes from negative to positive. Even though \( \Omega_e \) is increased further, the slope of \( n_2 \) remains positive.

5. Conclusion

In this paper, we have studied the self–Kerr nonlinear coefficient \( n_2 \) of the N-type system in detail. Since the probe field interacts with the middle two states \( |2\rangle \) and \( |3\rangle \), the slope and sign of \( n_2 \) near the resonance experience two dramatic changes between positive and negative with a unidirectional increase in the switching laser power, which shows the
reversible property of the system. Dressed state analysis has been presented to show the experimental result clearly.

Acknowledgments

This work was supported by the National Natural Science Foundation of China (Grant Nos. 11274112, 91321101, 11474092) the Fundamental Research Funds for the Central Universities (Grant No. WM1313003) and the key project of Shanghai Municipal Education Commission (Grant No. 14ZZ056).

References

[1] Courtois J-Y, Courty J-M and Reynaud S 1995 Quantum nondemolition measurements using a crossed Kerr effect between atomic and light fields Phys. Rev. A 52 1507
[2] Haas H A, Watanabe K and Yamamoto Y 1989 Quantum-nondemolition measurement of optical solitons J. Opt. Soc. Am. B 6 1138–48
[3] You H and Franson J D 2012 Theoretical comparison of quantum Zeno gates and logic gates based on the cross-Kerr nonlinearity Quantum Inf. Process 11 1627–51
[4] Li C Y, Zhang Z R and Sun S H 2013 Logic-qubit controlled-NOT gate of decoherence-free subspace with nonlinear quantum optics J. Opt. Soc. Am. B 30 1872–7
[5] Ferretti S and Gerace D 2012 Single-photon nonlinear optics with Kerr-type nanostructured materials Phys. Rev. Lett. 85 033303
[6] Chiang K S and Sammut R A 1993 Effective-index method for spatial solitons in planar waveguides with Kerr-type nonlinearity J. Opt. Soc. Am. B 10 704–8
[7] Towers I and Malomed B A 2002 Stable (2+1)-dimensional solitons in a layered medium with sign-alternating Kerr nonlinearity J. Opt. Soc. Am. B 19 537–43
[8] Matsko A B, Novikova I, Welch G R and Zubaityr M S 2003 Enhancement of Kerr nonlinearity by multiphoton coherence Opt. Lett. 28 96–98
[9] Niu Y P and Gong S Q 2002 Enhancing Kerr nonlinearity via spontaneously generated coherence Phys. Rev. A 73 053811
[10] Yang X D, Li S J, Zhang C H and Wang H 2009 Enhanced cross-Kerr nonlinearity via electromagnetically induced transparency in a four-level tripod atomic system J. Opt. Soc. Am. B 26 1423–34
[11] Kou J, Wang G R, Kang Z H and Gao J Y 2010 EIT-assisted large cross-Kerr nonlinearity in a four-level inverted-Y atomic system J. Opt. Soc. Am. B 27 2035–9
[12] Khao D X, Doai L V, Son D H and Bang N H 2010 EIT-assisted large cross-Kerr nonlinearity in a four-level inverted-Y atomic system J. Opt. Soc. Am. B 27 2035–9
[13] Field J E, Hahn K H and Harris S E 1991 Observation of electromagnetically induced transparency in collisionally broadened lead vapor Phys. Rev. Lett. 67 3062
[14] Harris S E 1997 Electromagnetically induced transparency Phys. Today 50 36–42
[15] Dawes A M C, Illing L, Clark S M and Gauthier D J 2005 All-optical switching in rubidium vapor Science 308 672–4
[16] Caglioti E, Trillo S, Wabnitz S and Stegeman G I 1988 Limitations to all-optical switching using nonlinear couplers in the presence of linear and nonlinear absorption and saturation J. Opt. Soc. Am. B 5 472–82
[17] Tran P 1999 All-optical switching with a nonlinear chiral photonic bandgap structure J. Opt. Soc. Am. B 16 70–73
[18] Milburn G J 1989 Quantum optical Fredkin gate Phys. Rev. Lett. 62 18
[19] Normandin R, Houghton D C, Simard-Normandin M and Zhang Y 1988 All-optical silicon-based integrated fiber-optic modulator, logic gate, and self-limiter Can. J. Phys. 66 833–40
[20] Sheng J T, Yang X H, Wu H B and Xiao M 2011 Modified self–Kerr nonlinearity in a four-level N-type atomic system Phys. Rev. A 84 053820
[21] Sheng J T, Miri M A, Christodoulides D N and Xiao M 2013 PT-symmetric optical potentials in a coherent atomic medium Phys. Rev. A 88 059904
[22] Yang K, Niu Y P, Chen D J and Gong S Q 2014 Laser frequency offset locking via tripod-type electromagnetically induced transparency Appl. Optics 53 2632–7
[23] Yang K, Niu Y P, Chen D J and Gong S Q 2014 Cavity linewidth narrowing by optical pumping-assisted electromagnetically induced transparency in V-type rubidium at room temperature J. Mod. Opt. 61 322–7
[24] Yang K, Niu Y P, Chen D J and Gong S Q 2014 Realization of cavity linewidth narrowing via interacting dark resonances in a tripod-type electromagnetically induced transparency system J. Opt. Soc. Am. B 31 144–8
[25] Wang H, Gooskey D and Xiao M 2001 Enhanced Kerr nonlinearity via atomic Coherence in a three-level atomic system Phys. Rev. Lett. 87 073601
[26] Wang H, Gooskey D and Xiao M 2002 Atomic coherence induced Kerr nonlinearity enhancement in Rb vapour J. Mod. Opt. 49 335–47
[27] Niu Y P, Gong S Q, Li R X, Xu Z Z and Liang X Y 2005 Giant Kerr nonlinearity induced by interacting dark resonances Opt. Lett. 30 3371–3
[28] Hamed H R, Radmehr A and Sahrai M 2014 Manipulation of Goos-Hanchen shifts in the atomic configuration of Mercury via interacting dark-state resonances Phys. Rev. A 90 053836