Optical preparation and measurement of atomic coherence at gigahertz bandwidth

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Abstract

We detail a method for the preparation of atomic coherence in a high-density atomic vapour of 87Rb, utilizing a coherent preparation scheme of off-resonant gigahertz bandwidth pulses. The scheme is found to be faster and more effective than techniques based on resonant interaction, such as coherent population trapping and population inversion. A numerical simulation of the preparation scheme is developed, and its efficiency in preparing coherent states is found to be close to unity at the entrance to the medium. The medium is then probed non-invasively with a laser field, the polarization of which is dependent upon the relative phase of the atomic coherence produced by the preparation fields.

The preparation of atomic coherence continues to draw a lot of attention [1–3], with potential applications in quantum computation and quantum information protocols [4]. Quantum memory [5] and entanglement of macroscopic systems [6], both of which are essential for large-scale quantum computers, have been achieved in high-density alkali metal vapours. Storage of light in such media has also been realized [7], allowing the storage and retrieval of photons at high speeds (gigahertz bandwidth) for optical quantum memory [8, 9]. This is advantageous for quantum computation, which calls for high-rate operations such as fast quantum gates based on Rydberg atoms [10, 11]. The technique we outline in this paper has the advantages of strong atom–light coupling to a high-density medium and the high-speed operation of pulsed light. Preparation of the medium is achieved using STImulated Raman Adiabatic Passage (STIRAP), a well-known method for the highly efficient transfer of population between two metastable states, facilitated via a stimulated two-photon transition involving an unstable intermediate state [12]. STIRAP has been generalized by allowing the three states to be degenerate manifolds [13]. In this case adiabatic transfer for any arbitrary pure or mixed initial ground state is possible. In this paper we describe theoretically an optical pumping technique which transfers atoms out of one of the dressed ground states; the remaining atoms which belong to the other dressed state are left in a superposition of bare states and are therefore part of a ground-state coherence. The phase and magnitude of this ground-state coherence is determined by the fields used to dress the atoms. These properties are measured via interaction between the atomic ensemble and the polarization of a probe field. Our technique thus provides a source of coherently prepared atoms suitable for quantum information processing, along with a means to non-invasively probe the atoms’ coherent state.

The remainder of this paper is as follows: in section 1 we describe the atom–light system necessary for the preparation scheme and provide details of the numerical model used in the simulation; in section 2 we show the result of the simulation; section 3 details a model of polarization rotation of a weak probe, the results of which are given in section 4. Finally, we demonstrate the optimal conditions for the STIRAP process in section 5, before drawing our conclusions in section 6.

1. Atom–light system

We consider the energy level scheme shown in figure 1(a). This consists of |g⟩, a Jg = 1 ground state which is a manifold of degenerate magnetic sublevels (|−⟩, |0⟩, |+⟩), an intermediate state |f⟩ with Jf = 0, and an excited state |e⟩. This three-level system is referred to as a ladder or cascade system, in which the levels successively increase in energy. For simplicity we do not consider the degeneracy of the excited state. The excited state population decays at a rate Γe, a fraction of which reaches the intermediate state. Similarly, the intermediate state population decays at a rate Γf, where it is distributed equally among the ground state sublevels.
As is typical for the STIRAP process, the excited state population is relatively long-lived. The ground–intermediate state coupling is via the pump electric field \( \mathbf{E}_p \), with associated slowly varying envelope \( \tilde{\mathbf{E}}_p = \tilde{E}_p e^{i \Delta t} + \tilde{E}_p e^{-i \Delta t} \). Here we have written the polarization state in the helical basis, where the components \( \tilde{E}_p \) and \( \tilde{E}_p \) stimulate the \( | - \rangle \leftrightarrow | f \rangle \) and \( | + \rangle \leftrightarrow | f \rangle \) transitions, respectively. The Faraday field \( \mathbf{E}_F \) stimulates the same transitions as the pump, but is of much lower intensity. Note that the remaining ground-state sublevel \( | 0 \rangle \) is only coupled to the other states via incoherent decay processes. Intermediate–excited state coupling is via the Stokes field \( \mathbf{E}_S \). The strength of a particular pump/Faraday-mediated transition is defined via the Rabi angular frequency \( \Omega_{\pm} = c m_d \mathbf{d}_f \cdot \mathbf{E}_S / \hbar \). Here \( \mathbf{d}_f \) is a reduced dipole matrix element describing the dipole coupling strength of a particular \( | L \rangle \rightarrow | L' \rangle \) transition; the coefficients \( c_m \) are factors governing the strength of \( | J, m_J \rangle \rightarrow | J', m_J' \rangle \) transitions. Likewise, the Stokes-mediated transition has a Rabi angular frequency \( \Omega_S = c_m \mathbf{d}_f \cdot \mathbf{E}_S / \hbar \). The pump (Faraday) fields are detuned from resonance by \( \Delta \rho (F) = \omega \rho (F) - \omega_0 \); the Stokes field is detuned by resonance by \( \Delta \rho = \omega_0 - \omega_0 \). The pulse sequence for the experiment is shown in figure 1(b).

Having described how energy levels are linked via the light fields, we first examine the effect of quantum interference [17], in the system. The reduced dipole matrix element of the pump transition \( \mathbf{d}_f \) is 5.177 \( \text{cm}^3 \text{GHz}^{-1} \) [20], with transition coefficients \( c_m = \frac{1}{4} \) for the three transitions \( | - \rangle \rightarrow | 0 \rangle, | + \rangle \rightarrow | f \rangle \); for the Stokes transition the reduced dipole element and transition coefficient are \( \mathbf{d}_e = 1.262 \text{cm}^3 \) [21] and \( c_m = - \frac{1}{\sqrt{2} m} \) (here \( m \) is the magnetic quantum number of an electron, \( m_0 \) is the Bohr radius). The Gaussian pulses have a full-width at half-maximum (FWHM) \( \delta t = 1 \text{ ns} \), with a 15\( \pi \) area for both components of the pump pulse and the Stokes pulse. The pump is initially linearly polarized at \( -\pi / 4 \) rad to the \( x \)-axis, and has a detuning \( \Delta \rho = 2 \pi (10 \text{ GHz}) \) from resonance. We assume two-photon resonance between the Stokes and pump fields, requiring that \( \Delta \rho + \Delta \rho = 0 \). The Faraday pulse has an area of \( 10^{-7} \pi \), is initially right-circularly polarized and is 5 GHz detuned from resonance. The medium has an atomic density of \( N_a = 10^{20} \text{ m}^{-3} \) (corresponding to a vapour temperature of \( \sim 150 ^\circ \text{C} \); figure 1 of reference [22] shows the nearly exponential dependence of number density on temperature).

2. Preparation of the medium

Before simulating the preparation stage of the atom–light interaction, we first examine the effect of quantum interference [23] in the system. The two competing paths to the excited state \( | - \rangle \rightarrow | f \rangle \rightarrow | e \rangle \) and \( | + \rangle \rightarrow | f \rangle \rightarrow | e \rangle \) lead to quantum interference: a well-known phenomenon in the interaction of multi-state systems with coherent light, see for example the review article [24]. It is instructive to transform the bare-atom set of basis states to a new set which takes into account the interaction with the light fields. We reformulate the ground state manifold \( \{| - \rangle, | 0 \rangle, | + \rangle \} \) using the Morris–Shore (MS) transformation [25] to \( \{| b \rangle, | 0 \rangle, | d \rangle \} \). The new basis states (the so-called dressed-atom states [15])

\[
| b \rangle = \frac{1}{\sqrt{\Omega}} (| - \rangle + | + \rangle),
\]

where the curly brackets \{ \} denote the anticommutator, and \( \sigma \) is the lowering operator. This master equation needs to be solved simultaneously with the slowly varying envelope form of the Maxwell wave equation

\[
\left( \frac{\partial}{\partial t} + \frac{c}{\partial z} \right) \tilde{E}_\pm = i \omega \rho (F) \frac{\mathbf{N}_a c m_d \mathbf{d}_f}{\epsilon_0} \tilde{\rho}_{ef},
\]

where (3) and (4) apply to the pump (Faraday) and Stokes pulse, respectively. The set of coupled equations (2)–(4) form the well-known Maxwell–Bloch equations (see for example [16, 17]), and are solved numerically using a Chebyshev pseudospectral time-domain method [18, 19].

To give an explicit example, we choose to model the \( \text{S}_3/2 (F = 1) \rightarrow \text{P}_3/2 (F = 0) \rightarrow \text{D}_3/2 \) transition found in \( ^{87}\text{Rb} \). This system has an intermediate-state decay rate of \( \Gamma_f = 2 \pi (6.065 \text{ MHz}) \), all of the atoms decaying out of state \( | f \rangle \) end up in the ground state. The excited state decays at the rate \( \Gamma_e = 2 \pi (0.66 \text{ MHz}) \); only a fraction (0.65) of the population decaying from \( | e \rangle \) ends up in \( | f \rangle \), the remaining fraction decays to other states not included in our five-level system. The reduced dipole matrix element of the pump transition \( \mathbf{d}_f \) is 5.177 \( \text{cm}^3 \text{GHz}^{-1} \) [20], with transition coefficients \( c_m = \frac{1}{4} \) for the three transitions \( | - \rangle \rightarrow | 0 \rangle, | + \rangle \rightarrow | f \rangle \); for the Stokes transition the reduced dipole element and transition coefficient are \( \mathbf{d}_e = 1.262 \text{cm}^3 \) [21] and \( c_m = - \frac{1}{\sqrt{2} m} \) (here \( m \) is the magnetic quantum number of an electron, \( m_0 \) is the Bohr radius). The Gaussian pulses have a full-width at half-maximum (FWHM) \( \delta t = 1 \text{ ns} \), with a 15\( \pi \) area for both components of the pump pulse and the Stokes pulse. The pump is initially linearly polarized at \( -\pi / 4 \) rad to the \( x \)-axis, and has a detuning \( \Delta \rho = 2 \pi (10 \text{ GHz}) \) from resonance. We assume two-photon resonance between the Stokes and pump fields, requiring that \( \Delta \rho + \Delta \rho = 0 \). The Faraday pulse has an area of \( 10^{-7} \pi \), is initially right-circularly polarized and is 5 GHz detuned from resonance. The medium has an atomic density of \( N_a = 10^{20} \text{ m}^{-3} \) (corresponding to a vapour temperature of \( \sim 150 ^\circ \text{C} \); figure 1 of reference [22] shows the nearly exponential dependence of number density on temperature).
\[ |d\rangle = \frac{1}{\Omega_p}(\Omega_-|\rangle - \Omega_+|+\rangle), \tag{6} \]

are, respectively, coupled and uncoupled from the state |f\rangle; the magnetic sublevel |0\rangle remains uncoupled. For simplicity, we take the polarization state of the pump to be fixed throughout the experiment, and thus the dressed states also remain fixed.

The Hamiltonian of the transformed system

\[ \mathcal{H}_{\text{RWA}}' = -\frac{i}{\hbar}(\Omega_p|f\rangle\langle b| + \Omega_S|e\rangle\langle f| + 2\Delta_p|f\rangle\langle f| + 2(\Delta_S + \Delta_p)|e\rangle\langle e|) + \mathbf{H.c.} \tag{7} \]

shows that the transition |b\rangle \leftrightarrow |f\rangle is mediated with the Rabi angular frequency \[ \Omega_p = \sqrt{\Omega_+^2 + \Omega_-^2}. \]

This tells us that the atomic state |b\rangle is associated with the polarization state of the pump field \( \mathbf{E}_P \); similarly, the uncoupled state |d\rangle is associated with a field orthogonal to \( \mathbf{E}_P \), the magnitude of which is zero in the preparation stage. We discuss the implications of this later.

If we begin with an atomic ensemble in a mixed state, the initial ground state density operator in the MS transformed basis is

\[ \hat{\rho}_{\text{initial}} = \frac{1}{4}(|d\rangle\langle d| + |0\rangle\langle 0| + |b\rangle\langle b|), \]

i.e. the three possible states are evenly populated and there is no coherence amongst them. The effect of the preparation fields is to effect a two-photon transition between the states |b\rangle and |e\rangle. The final ground state density operator is then

\[ \hat{\rho}_{\text{final}} = \frac{1}{4}(|d\rangle\langle d| + |0\rangle\langle 0| + \delta|b\rangle\langle b|), \]

where \( \delta \to 0 \) for complete population transfer. Examining the form of the dressed-atom states in equations (5) and (6) we see that they are orthogonal and are coherent superpositions of the bare-atom states |−\rangle and |+\rangle. Therefore asymmetry in the populations of the dressed states leads to an increase in the coherence of the ground-state subsystem {−, +}. The aim of the preparation process is to create a partially coherent ground state and the efficiency of the preparation stage can be parameterized as the degree of coherence \( p_{+-} = |\rho_{+-}|/\sqrt{\rho_{++}\rho_{--}} \), which takes a value from zero (an incoherent mixture) to the maximum allowed value of unity (a pure state).

Figure 2 shows the theoretical results of pulse propagation in the preparation stage. In parts (a)–(d) the field envelopes are plotted. The peak of the Stokes pulse enters the medium at \( t = 0 \) ns, followed by the pump at \( t = 0.6 \) ns. This pulse spacing amounts to a time separation of one \( 1/e \) width, which is the optimal separation for STIRAP using Gaussian envelope functions [12]. The front of the Stokes pulse is seen to traverse the medium at close to the speed of light in vacuo. (c) The Stokes field magnitude versus time, at \( z = 0 \) (dashed curve) and \( z = 22.5 \) mm (solid curve); (d) shows the same information for the pump field.
media, the preparation fields experience less absorption/gain and as a consequence can travel greater distances before becoming distorted. The resultant media are prepared at high efficiency to a greater depth. For example, for a room temperature medium of atomic density $N_e = 10^{18} \text{ m}^{-3}$, the preparation efficiency is approximately unity for a distance of several metres. In our simulation, however, the use of a medium with a large atomic density ($N_e = 10^{20} \text{ m}^{-3}$) was necessary to observe large rotations during the measurement stage, as described in the next section.

3. A simple model of polarization rotation

After the STIRAP preparation stage we are left with a medium possessing a ground-state subsystem prepared to a high degree of coherence which is stable against decay (we ignore particle–particle interactions such as collisions [22] which may serve to dephase the partially coherent state). The rest of the atomic population is found in the metastable excited state from where it will eventually filter back to the ground state, but only over a timescale longer than the duration of the experiment. To a fair degree of accuracy we can then treat the medium as if it were in a stable state. In the measurement stage, we apply a weak probing field to the medium. For a weak enough probe beam the medium is unaffected by the passage of the field, which leaves us free to assume steady-state conditions.

Working in the frequency domain we have a polarization density of $\mathbf{P} = \frac{1}{\hbar} \chi \mathbf{E} = N_e \langle \hat{d} \rangle$, where $\chi$ is the susceptibility tensor. We can write the field and expectation value of the dipole operator as column vectors, giving the expression

$$\hat{\chi} \left( \mathbf{E}_+ \mathbf{E}_- \right) = \frac{2N_e d_0}{\hbar} \left( \hat{\rho}_{zz} - \hat{\rho}_{xx} \right).$$

Here the dipole matrix element $d_0$ is equal to the reduced dipole matrix element $d_{\text{red}}$ multiplied by $e_{nm}$, the relative coefficient of the $|\langle-\rangle, |\rangle\rangle \rightarrow |f\rangle$ transitions. Note that the magnitude of $e_{nm}$ is equal for both transitions (due to the symmetry of the electronic wavefunction [27]). The steady-state values of the coherence terms can be derived from the Bloch equations, and are found to be

$$\dot{\rho}_{zz}^\text{st} = \frac{i d_0}{2\hbar} \mathbf{E}_+ \mathbf{E}_- \rho_{zz}^\text{st} + \mathbf{E}_- \mathbf{E}_+ \rho_{zz}^\text{st}$$

and

$$\dot{\rho}_{zz}^\text{st} = \frac{i d_0}{2\hbar} \mathbf{E}_- \mathbf{E}_+ \rho_{zz}^\text{st}$$

assuming that the population of the intermediate state $\rho_{zz}^\text{st} \approx 0$. Substituting these steady-state solutions of the coherence into equation (8), we can express the susceptibility tensor as

$$\hat{\chi} = \frac{N_e d_0}{\hbar e_0} \frac{i}{\Gamma_f / 2 - i\Delta_F} \left( \rho_{zz}^\text{st} - \rho_{xx}^\text{st} \right) \left( \rho_{zz}^\text{st} - \rho_{xx}^\text{st} \right).$$

In the absence of a coherence between states $|\rangle$ and $|\rangle$, the susceptibility tensor is diagonal and thus the field components $\mathbf{E}_+$ and $\mathbf{E}_- \langle \rangle$ are independent of each other. However, in the presence of a coherence this is not the case and there will be interference between the two field components. The normal modes of the field, i.e. the field polarizations that propagate independently of each other, can be found by diagonalizing $\hat{\chi}$.

However, we noted in section 2 that during the STIRAP process the polarization state of the pump field defines the coupled dressed state $|b\rangle$ (and the orthogonality condition determines the uncoupled state $|d\rangle$). The fields associated with the states $|b\rangle$ and $|d\rangle$ are the normal modes of the medium.

A convenient visual representation of light polarization is the Poincaré sphere. Points in this three-dimensional space correspond to the column vector (here $T$ denotes the transpose operation) $S = (S_1 S_2 S_3)^T$, the components of which are, respectively, the intensity difference between linearly polarized light in the $x$ and $y$ directions, the intensity difference between linearly polarized light at an angle $+\pi/4$ and $-\pi/4$ rad to the $x$-axis, and the intensity difference between left and right circularly polarized light. Note that orthogonality is represented by antipodal points. A light field with temporarily and spatially varying polarization is generally described by a surface. The vector/surface is often normalized by the total light intensity, and for fully polarized light each point lies on a sphere of unit radius. The evolution of the polarization vector is implicit in the Maxwell–Bloch equations, but to aid the interpretation of the numerical solution to these equations, we note that the torque equation of motion provides a simple analogy of birefringence [14, 28]. The equation describes the spatial evolution of the polarization vector $\mathbf{S}$ in response to the anisotropy of the medium, represented by the birefringence vector $\mathbf{a}$:

$$\frac{d\mathbf{S}}{dz} = \mathbf{a} \times \mathbf{S}.$$  (12)

The geometric interpretation on the Poincaré sphere is that $\mathbf{a}$ provides the instantaneous rotation axis and rotary power for the evolution of $\mathbf{S}$. Note the limitations of this simple picture, however, in that it assumes monochromatic waves in a time-independent medium with zero losses.

The birefringence vector points in the direction of the preponderance of atoms in the $|d\rangle$ state. The Stokes parameters of the anisotropy vector can be related to the density matrix elements of the ground-state subsystem via the expression [29]

$$\mathbf{a} = (-2\text{Re} \left[ \hat{\rho}_{zz}^\text{st} \right], 2\text{Im} \left[ \hat{\rho}_{zz}^\text{st} \right], \hat{\rho}_{zz}^\text{st} - \rho_{xx}^\text{st})^T.$$  (13)

Note the third element in this vector which is due to an imbalance in the populations of the states $|\rangle$ and $|\rangle$. This is the cause of the traditional paramagnetic Faraday effect [30], and is a manifestation of circular birefringence. If the populations are equal and a ground-state coherence exists, the medium will be linearly birefringent, i.e. will respond differently to two orthogonal linearly polarized field components. In the general case the medium is elliptically birefringent.

In the case of the birefringence vector $\mathbf{a}$ being independent of distance inside a medium of uniform density, the degree to which the polarization vector rotates is proportional to both distance $z$ and atomic density $N_e$. To produce rotations of order $\pi$ rad over distances of a few millimetres it was found necessary to use a high-density medium of $\sim 10^{20}$ atoms per cubic metre. The same rotation can be achieved at lower densities concomitant with a proportionally longer medium.

We compare the results of the numerical simulation and torque equation in the next section.
In consideration of the complications of the STIRAP process seen in figures 2 and 3, a 2 mm long medium will be used to ensure a relatively homogeneous sample for the measurement stage. In this stage, the Faraday pulse interacts with the medium after the preparation fields have exited: the pulse has a detrimental effect on the measurement process: the pulse does not rotate as a single entity, rather the variation of dispersion over its bandwidth leads to differential rotation. Thus for each position inside the medium the Stokes vector varies in time. This manifests itself upon the Poincaré sphere as the spreading out from a single point. Figure 4(d) shows the field on the polarization ellipse. Here the polarization state corresponding to the peak of the pulse is shown at the entrance and exit of the medium, along with one of the normal modes of the medium (the other mode is orthogonal to this).

We have previously considered the regime in which the ground state is stable against dephasing mechanisms which have the tendency to lessen the degree of coherence of the atomic subsystem we are interested in. Dephasing can be the result of, for example, collisions between two atoms producing a relative phase between them, reducing the overall coherence of the entire atomic ensemble. To model the effects of dephasing we now add a 27 MHz dephasing term to the master equation (this rate is chosen so as to significantly affect the coherence during the few nanoseconds that is the duration of the simulation). Dephasing causes the sublevel coherence to decay exponentially towards an incoherent mixture, i.e. the degree of coherence $p_{+-}$ decreases with time. Thus the magnitude of the birefringence vector (see equation (13)) decreases, the effect of which is a gradual reduction in the rotary power of the birefringent medium with time, though the axis of rotation remains pointing in the same direction of the Poincaré sphere. In figures 4(b) and (e) we see the effect of dephasing, where due to reduction in the level of coherence the polarization of the Faraday pulse rotates to a lesser degree than the case where dephasing mechanisms were ignored.

Finally, we consider the effect of an energy difference between the two sublevels involved in the coherence. This simulates an applied longitudinal magnetic field used in Faraday rotation. The energy difference causes precession of the atomic spin, analogous to the Larmor precession of magnetic moments around an applied magnetic field [32]. The degeneracy of the levels is broken to such an extent that the partially coherent state precesses at a rate of 27 MHz. With the system we are modelling in this paper (see section 1) the corresponding applied magnetic field is around 20 Gauss: typical of a Faraday rotation experiment. Due to precession, by the time the Faraday pulse enters the medium, the birefringence vector has rotated to a new direction, as observed in figures 4(c) and (f). The duration of the light pulse is 1 ns, therefore the birefringence vector rotates a negligible amount in the time it takes the light to traverse the medium. Since the precession of the birefringence vector is in the $xy$-plane, it does not have a detrimental effect on the measurement process: the light polarization rotates around the birefringence vector in a
Figure 5. Effect of time-separation of the pump and Stokes pulses on the STIRAP process. (a) Population of the excited state upon completion of the preparation stage, shown against the separation time between the Stokes and pump pulses. (b) The maximum population seen in the intermediate state. The asymmetry around $t_{sep} = 0$ and the trough on the positive side is the signature of the STIRAP process. (c) Rotation angle $\theta$ of the Faraday pulse is shown (left axis, solid line) overlaid with the magnitude of the birefringence vector (right axis, data points, the number of which has been reduced for clarity).

It is well known that the overlap of the Stokes and pump envelopes plays a part in the efficiency of the STIRAP process [12, 33]. Figure 5(a) shows the population transferred to the excited state versus the separation time between the peaks of the incident Stokes and pump pulses. For positive separation, the Stokes pulse precedes the pump, which is the correct order for adiabatic population transfer. This can be seen from the trough in figure 5(b), where transfer to the intermediate state is at a minimum. Note that due to the pump beam being detuned off resonance, the intermediate state is only transiently populated during the preparation stage. The asymmetry is not mirrored in figure 5(a) because by carrying out the transfer faster than the decay rate we are less harshly punished for going on an excursion to the intermediate state. Figure 5(c) shows both the length of the birefringence vector $a$ and the rotation angle $\theta$ around this vector experienced by the peak of the Faraday pulse. The magnitude of the birefringence vector and the rotation angle are clearly linked.

6. Conclusion

We have demonstrated a theoretical method for the preparation and measurement of a coherence in the ground state of a high-density atomic medium, using gigahertz bandwidth pulses. The technique we have described has a number of advantages over other preparation schemes, such as those based upon spontaneous emission. These schemes, which include coherent population trapping (CPT) [34] and population inversion, take many excited state lifetimes to transfer the medium into its final state, which substantially limits the bandwidth of the operation. Also, the efficiency of CPT-based schemes decreases with increasing atomic density, as photons scatter from the pumping field and continue to interact with the sample [35]. For these reasons, transfer via an off-resonant coherent mechanism such as the one described in this paper is preferred, in which there is a certain amount of freedom to choose the final state of the atom, and the timescales over which the process is carried out are shorter than those necessary for incoherent pumping. In addition, despite the use of a high-density medium, which was necessary due to the requirement for strong coupling to the probe field, we have shown it is possible to produce a uniformly prepared sample of a much greater length than that possible with resonant processes. We have shown that the sample can be probed non-invasively with a gigahertz bandwidth pulse.

The parameters of the atom–light system were chosen to model the $5S_{1/2}(F = 1) \rightarrow 5P_{3/2}(F = 0) \rightarrow 5D_{5/2}$ transition found in $^{87}$Rb, therefore our technique is readily amenable to experimental investigation.

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