Causes of EIT Intensity Correlation Power Broadening

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(Dated: January 5, 2015)

Abstract

EIT noise correlation spectroscopy holds promise as a simple, robust method for performing high resolution spectroscopy used in optical magnetometry and clocks. Of relevance to these applications, we report here on a measurement of and a theory model for power broadening of EIT noise correlation resonances. Taken together they identify physical antecedents of noise correlation power broadening.

PACS numbers: 42.50.Gy, 32.70.Jz, 34.80.Pa
I. INTRODUCTION

When laser fields interact with atomic vapors to induce atomic coherence between dipole-forbidden quantum states, such as between two ground states in a Λ-type three-level system, the optical properties of the atomic vapor are dramatically altered in controllable and useful ways. Simultaneously, interactions with the atomic vapor modifies the participating light fields. A well-known representative coherent phenomenon is electromagnetically induced transparency (EIT) [1], which owes its spectrally narrow transmission window and ultra-steep dispersion to optical pumping into a dark state, which is understood as quantum interference between transition pathways. The controllable optical properties that accompany EIT and related coherent phenomenon such as coherent population trapping (CPT) and electrically induced absorption (EIA) are particularly attractive for applications [2] such as atomic clocks [3], magnetometry [4], quantum information and communication schemes [5, 6] and quantum computation [7]. This widespread interest is motivation for understanding properties of the light fields emerging from interaction with the atomic coherence.

Over the past ten years experiments have ultimately led to significant insight into the intensity fluctuations of light emerging from coherently prepared media such as EIT in atomic vapor. Before entering such media inherent optical phase noise is typically much more spectroscopically relevant than residual intensity noise. Although phase noise is not easily detectable by photodetectors, light-matter interactions like typical atomic absorption or EIT map the phase noise onto the light’s amplitude (and thus, intensity) fluctuations. The phase modulation to amplitude modulation (PM → AM) conversion processes [8, 9] can lead to excess noise that is usually metrologically undesirable. However, PM → AM conversion processes impart the noise with spectroscopic information making it a useful spectroscopic tool [10–12]. Furthermore, monitoring correlations between the coherence-derived light fluctuations in the two EIT (or EIA) laser fields, a.k.a. EIT noise correlation spectroscopy, shows great promise for a range of applications due to a narrow, power-broadening resistant resonance described below and studied previously in [13–16].

Intensity fluctuations of the two EIT laser fields become strongly correlated on EIT resonance, and abruptly strongly anti-correlated slightly off-resonance. This correlation-anticorrelation switching near EIT gives rise to a narrow correlation resonance, whose linewidth is up to an order of magnitude more narrow than the underlying EIT trans-
parenity peak, and has been shown in some cases to be more narrow that the zero-power EIT transit-limited lifetime [14]. The abruptness of the correlation switch arises because the light fluctuations directly probe the coherences causing the EIT [15], typically obscured at high powers by the sensitivity-defeating consequences of power-broadening [17].

This narrow correlation resonance is of further scientific and technological interest because it has been shown in experiment and in theory models to be power-broadening resistant at laser intensities which broaden the underlying EIT resonance [13, 17]. Coherence-based applications that exploit this correlation resonance will have an advantage in sensitivity over those with the usual optical schemes, and for many applications, the prospect of working at higher optical powers without broadening is clearly technologically attractive. More recent experiments indicate that the noise correlation resonance does eventually broaden at significantly higher laser intensities, but questions remain about the physical cause(s) of the broadening.

In this paper, we identify physical processes that contribute to the slope of this observed EIT noise correlation power broadening. We report on results from new experiments and a simple quantitative quantum optics theory model applied to intensity fluctuations of orthogonally polarized fields in the EIT Hanle configuration of the $^{87}\text{Rb} \, D_1$ line. We have studied the dependence of the intensity cross-correlation noise statistic, referred to as $g^2(0)$, on magnetic field gradients across the atomic vapor cell, beam intensity, beam size, one-photon detuning, residual amplitude modulation (RAM) of the beam, cell temperature, and buffer gas pressure. In Section III we report in detail on a few of these effects starting with new experiments and findings that probe EIT noise correlation power-broadening. Section IV is a brief description of the theory model and its connection to the experimental results, followed by a summary in Section V.

II. EXPERIMENTAL SETUP

A schematic of our experimental setup is shown in Fig. II. We induce EIT in a warm $^{87}\text{Rb}$ buffer gas vapor cell using the Hanle configuration, which employs orthogonal circularly polarized components $\sigma+$ and $\sigma-$ of a single linearly polarized laser to induce an atomic coherence between degenerate ground state Zeeman sublevels. The vapor cell is heated to 50$^\circ$C, and contains a buffer gas of either 2 Torr Ne or 10 Torr (5 Torr Ne + 5 Torr Ar)
with a cell length of $l = 5$ cm or 8 cm, respectively. In order to observe Zeeman EIT, the atomic vapor cell is shielded from the ambient magnetic field of the laboratory with three nested layers of $\mu$-metal shielding. A nearly homogeneous magnetic field is created by a solenoid inside the innermost shield layer. This longitudinal B-field is co-linear with the beam’s propagation direction and is used to split the ground state Zeeman sublevels, and thus to control the two-photon detuning, $\Delta$. A small coil mounted far from the cell is used in some experiments to introduce a magnetic field gradient across the vapor cell.

We tune a commodity 795 nm non-narrow linewidth ($\sim 55$ MHz) free-running diode laser using temperature and current control to the $F=2 \rightarrow F'=1$ hyperfine transition of the $^{87}\text{Rb}$ D1 line. To ensure stability of the laser system against long-term temperature drifts in
the lab environment, the center of the laser line is locked to the transition using a dichroic atomic vapor laser lock (DAVLL) \[18\], and a saturated absorption setup is used to monitor the laser wavelength. The large spectral bandwidth of such a “noisy” laser is desirable for EIT noise correlation studies because its large intrinsic phase noise is much more relevant spectroscopically than the small remaining laser intensity instability \[19\]. We have measured the intrinsic laser intensity noise, or RAM, for our free-running laser diode to be roughly $\sim 0.2\%$. This inexpensive diode laser was chosen for compatibility with device applications such as CPT atomic clocks \[20\] and magnetometers \[21\] and to demonstrate these simple lasers can provide a robust solution for in-the-field atomic-optical applications where cavity alignment and stability are limited by the environment \[22\].

The $\sigma^+$ and $\sigma^-$ components of the linearly polarized laser field couple two degenerate ground state Zeeman sublevels to a common upper level, creating a $\Lambda$-type three-level system. The resulting EIT is observed as a peak in the transmission of light as the atoms are subject to a magnetic field that is swept through $B = 0$.

After the atomic vapor cell, a quarter waveplate and polarizing cube beamsplitter separate the $\sigma^+$ and $\sigma^-$ channels of the transmitted light, and each are individually detected with amplified photodiodes, whose bandwidths are limited to $<1$ MHz. For each applied magnetic field, or two-photon detuning, the intensity in each polarization channel is recorded at a rate of 1 MHz for 4 ms with enough dynamic range to detect the AC part, $\delta I_{\pm}$, of the photodiode signal, that is, the signal with the mean value subtracted.

Of particular interest is the degree of correlation of the two fluctuating light fields, measured by the normalized intensity cross-correlation statistic, $g^2(0)$:

$$g^2(0) = \frac{\langle (\delta I_+)(\delta I_-) \rangle}{\sqrt{\langle (\delta I_+)^2 \rangle \langle (\delta I_-)^2 \rangle}}. \tag{1}$$

The numerator of $g^2(0)$ averages the products of the AC parts of the two signals, and the denominator normalizes the result such that perfect correlation outputs $g^2(0) = +1$ and perfect anticorrelation yields $g^2(0) = -1$. The utility of this intensity cross-correlation statistic in probing EIT resonances was first reported in \[13\].

Fig. 2(a) shows a representative $g^2(0)$ dataset, which has a dual structure consisting of an outer envelope that can approach complete anticorrelation, $g^2(0) = -1$, near zero-detuning and an inner central resonance peak arising from correlation switching to perfect correlation at zero two-photon detuning, $\Delta = 0$. The superimposed line in Fig. 2(a) is
FIG. 2. (color online) (a) A representative $g^2(0)$ EIT intensity correlation curve. Data points are the average of $g^2(0)$ values calculated from 10 consecutive data acquisition periods, and the (red) curve is from our theory model. (b) The central correlation peak, with fitted Gaussian curve to determine FWHM. These data were taken in a 10 Torr buffer gas cell with beam intensity $I = 0.7$ mW/cm².

from our theory model (in Section IV below) and is in good agreement with the experiment. Fig. 2(b) highlights the narrow correlation resonance feature, the linewidth of which we experimentally determine using the full width at half max (FWHM) of a fitted Gaussian curve.

III. EXPERIMENTAL RESULTS & DISCUSSION

A series of experiments were performed to investigate the relationship of the central correlation peak’s linewidth to laser power. Fig. 3(a) shows the central correlation peak for different laser powers, and Fig. 3(b) summarizes the central peak’s FWHM growth with laser intensity. We observe that beyond a low power threshold, in this case $I > 1$ mW/cm²,
the correlation linewidth broadens linearly with laser power. We have defined the “broadening rate” as the slope of this linear trend. For the experimental conditions in Fig. 3, the broadening rate is 0.85 kHz per mW/cm\(^2\).

Fig. 4(a) shows a series of correlation peak power broadening experiments using four different beam sizes (radii \(r = 0.25 - 0.5\) cm). With correlation peak FWHM plotted as a function of intensity (total input power normalized by beam area), we learn that surprisingly, beam size has little, if any, affect on the broadening rate. Keeping intensity fixed, we note that experiment and theory for these buffer gas cells do indicate a very weak dependence on the beam size, though it is more relevant for the threshold power at which correlation peak power broadening first sets in. Fig. 4(b) plots the noise correlation peak FWHM as a function of intensity for two buffer gas vapor cells, 2 Torr and 10 Torr, using similar beam sizes. Here we see a clear dependence of the broadening rate, with the 10 Torr cell exhibiting a noise correlation broadening rate four times that of the 2 Torr cell.

We use the following equation from [23] to calculate the loss of coherence due to diffusion of the atoms out of the beam, which contributes directly to the ground state decoherence:

\[
\Gamma_{\text{dif}} = \frac{d_o p}{P} \left[ \left(\frac{2.41}{r}\right)^2 + \left(\frac{\pi}{7}\right)^2 \right],
\]  

(2)

where \(\Gamma_{\text{dif}}\) is the rate at which the atoms diffuse out of the beam in Hz, \(d_o\) is the diffusion constant in cm\(^2\)/s, \(P\) is the buffer gas pressure in Torr, and \(r\) and \(l\) are the beam radius and cell length of the cell in cm. With \(d_o = 0.23\) cm\(^2\)/s for Ne and 0.16 cm\(^2\)/s for Ar [24], we calculate that an increase in the beam size from \(r = 0.25\) cm to 0.5 cm, in either the 2 Torr or 10 Torr cell, causes the diffusion rate to be reduced by a factor of 4.0. Our results in Fig. 4(a) indicate that this change in diffusion rate does not correspond to an increased power broadening rate. Similarly, when beam radius is kept constant, the diffusion rate is reduced by a factor of 6.0 times for the longer 10 Torr mixed species cell compared with the shorter 2 Torr Ne cell. In view of the results in Fig. 4(a), it is unlikely that this modest improvement of the diffusion rate/transit time between the 2 and 10 Torr cells accounts for the slope discrepancy seen in Fig. 4(b).

We understand the increased power broadening rate for the higher pressure cell in Fig. 4(b) as a consequence of increased collisional broadening of the excited state. Collisional broadening causes the width of the excited state to broaden at a rate of 10 MHz/Torr for Ne and 18 MHz/Torr for Ar [25], yielding a excited state in the 10 Torr cell that is broadened
FIG. 3. (color online) (a) The $g^2(0)$ central correlation peak for four beam intensities. (b) FWHM of the central correlation peak for a series of beam intensities. Note that the FWHM of the peaks in (a) are indicated in (b) by corresponding open datamarkers. These data were taken in a buffer gas pressure of 10 Torr.
FIG. 4. (color online) (a) FWHM of the central correlation peak as a function of intensity for four beam sizes, radii $r = 0.25 - 0.5$ cm, in a 10 Torr buffer gas vapor cell. (b) FWHM of the central correlation peak as a function of intensity in two buffer gas vapor cells, 2 Torr (open circles) and 10 Torr (filled squares). See text for vapor cell details.

to more than two times the linewidth of our laser, and more than five times as in the 2 Torr cell. Rubidium atoms in the higher pressure buffer gas are resonant with the laser more often. For signal averaging this increased time on resonance corresponds to a larger effective Rabi frequency; thus increased buffer gas pressure leads to larger power broadening slopes.

We also observe (data not shown) that the correlation resonance broadening rate does not appreciably change over the range of temperatures for which there is an unmistakable central feature in the $g^2(0)$ curve, $40 - 60^\circ$C in our studies. Below $40^\circ$C the atomic vapor is not dense enough to produce a discernible AC signal, and above $60^\circ$C, absorption effects dominate and obscure the correlation effects. Our studies in this temperature range indicate that radiation trapping is not an underlying cause of the correlation peak power broadening.

In an effort to understand the physics underlying the EIT noise correlation broadening rate and the possible role of spatial inhomogeneities in the two-photon detuning across the cell, we controllably introduced a $B$-field gradient within the $\mu$-metal shielding using a small current-carrying coil. The coil was placed with its axis parallel to the optical axis. A fluxgate magnetometer was used to map the longitudinal component of the coil’s $B$-field gradient within the magnetic shielding for a range of currents through the coil. In the region of the
FIG. 5. (color online) For each applied $B$-field gradient value, the correlation peak FWHM was measured as a function of intensity, and the broadening rate was determined by a linear fit. For reference, the 0 mG/cm point was determined by the data in Fig. 3(b).

atomic vapor cell, the gradient was found to be approximately spatially constant and linear in applied current, as expected.

For each applied $B$-field gradient value, the correlation peak FWHM was measured as a function of intensity, and the broadening rate was determined by a linear fit, as in Fig. 3(b). Fig. 5 plots the broadening rate of the central correlation peak versus applied $B$-field gradient. The data suggests that the broadening rate grows linearly with applied $B$-field gradient. Interpretation of this result is given in Section V.
IV. THEORY MODEL

A straightforward quantum optics theoretical model quantitatively captures the behavior of EIT-derived intensity correlations. In order to establish a theory baseline, the simplest model consists of a standard a three-level Λ-system with ground states |1⟩ and |2⟩, and excited state |0⟩. Referring to Fig. 1 let γ₁ be the bare excited state population decay rate, γ₂ the depolarization of the dipole coherences, Ω± the Rabi fields, D their common one-photon detuning, Δ the two-photon detuning, Γ the bare ground state decoherence rate, Γ₁ the ground state population redistribution rate, and d = ρ₁₁ − ρ₂₂ the ground state population difference. The density matrix equations are well known, but reproduced here for completeness and to fix our notation. They read:

\[ \partial_t \rho_{00} = -\gamma_1 \rho_{00} + \frac{i\Omega_+}{2}\rho_{01} + \frac{i\Omega_-}{2}\rho_{02} + \text{c.c.} \]  
\[ \partial_t \rho_{11} = \frac{\gamma_1}{2}\rho_{00} - \Gamma_1(\rho_{11} - \rho_{22}) - \frac{i\Omega_+}{2}\rho_{01} + \text{c.c.} \]  
\[ \partial_t \rho_{22} = \frac{\gamma_1}{2}\rho_{00} - \Gamma_1(\rho_{22} - \rho_{11}) - \frac{i\Omega_-}{2}\rho_{02} + \text{c.c.} \]  
\[ \partial_t \rho_{01} = -(\gamma_2 - i\delta_1)\rho_{01} + \frac{i\Omega^*_+}{2}(\rho_{00} - \rho_{11}) - \frac{i\Omega^*_+}{2}\rho_{21} \]  
\[ \partial_t \rho_{02} = -(\gamma_2 - i\delta_2)\rho_{02} + \frac{i\Omega^*_+}{2}(\rho_{00} - \rho_{22}) - \frac{i\Omega^*_+}{2}\rho_{21}^* \]  
\[ \partial_t \rho_{12} = -(\Gamma - i(\delta_1 - \delta_2))\rho_{12} - \frac{i\Omega_+}{2}\rho_{02} + \frac{i\Omega^*_+}{2}\rho_{01}^* \]

Here Δ = δ₁ − δ₂ is the two photon detuning, and thus δ₁,₂ = D ± Δ/2. For the cross-correlation statistic in the optically thin cell limit, the intensity fluctuations of each polarization, δI±, are proportional to the fluctuations in the associated dipole coherences, δI± ∼ Ω±*δρ₀₁, Ω±*δρ₀₂.

Our main theory assertion is that a careful averaging of the steady state response of equations Eqs. 3-8 over an ensemble of one-photon detunings quantitatively describes the observed phenomena in the noise correlation spectroscopy experiments. For typical
commodity laser diodes, the ensemble consists of one-photon detunings associated with laser phase noise. Assuming that the laser phase noise is spectrally wider than the atomic one-photon transition, $\gamma_1$, we take the ensemble of one-photon detunings to be a flat distribution. Additionally our numerical model accommodates intrinsic RAM as well, either correlated or uncorrelated with phase noise.

Instead of solving Eqs. (3)-(8) exactly in the time domain with temporally changing detunings, we separate ‘fast’ and ‘slow’ timescales, with respect to $\gamma_1$, of the response and ensemble average. Among ‘fast’ processes are the laser phase noise and the optical response of the dipole coherences $\rho_{01}$ and $\rho_{02}$. ‘Slow’ processes include the evolution of the ground state coherences, $\rho_{12}$, and the ground state populations, $d$, since they evolve at a rate slower than the bare optical pumping rate in these vapor cells. All fast processes are averaged over the ensemble of one-photon detunings, $D$, whereas the slow processes depend on the ensemble centroid $D_0 = \langle D \rangle$ only. Throughout, in keeping with the notation in Eq. (1), brackets $\langle \rangle$ represent averaging over the ensemble.

In practice, solving Eq. (6) and Eq. (7) for the steady state values of $\rho_{01}$ and $\rho_{02}$ and substituting those into Eqs. (8), (4) and Eq. (5) gives,

$$\tilde{\Gamma}_{\rho_{12}} = \frac{-\Omega_+ \Omega_+^* \gamma_2}{4(\gamma_2^2 + D_0^2)} + \frac{i\Omega_+ \Omega_+^* D_0 d}{4(\gamma_2^2 + D_0^2)}$$

(9)

$$\tilde{\Gamma}_1 d = \frac{(|\Omega_-|^2 - |\Omega_+|^2)\gamma_2}{8(\gamma_2^2 + D_0^2)} + \frac{D_0(i\Omega_+ \Omega_- \rho_{12} + c.c.)}{4(\gamma_2^2 + D_0^2)}$$

(10)

where $\tilde{\Gamma} = \Gamma - i\Delta' = \Gamma - i\Delta + \frac{|\Omega_+|^2}{4(\gamma_2 - iD_0)} + \frac{|\Omega_-|^2}{4(\gamma_2 + iD_0)}$ and $\tilde{\Gamma}_1 = \Gamma_1 + \frac{(|\Omega_+|^2 + |\Omega_-|^2)\gamma_2}{8(\gamma_2^2 + D_0^2)}$. That is, $\Gamma'$ and $\Delta'$ are real. We have ignored the excited state population as negligible but have included terms of order the excited state population divided by the excited state lifetime (thus accounting for leading order power broadening in the one-photon transitions.)

Since the $\rho_{01}$ and $\rho_{02}$ evolve on much shorter time scales they respond to a much larger band of laser frequencies inside the phase noise bandwidth, particularly in vapor cells with high pressure buffer gas or a quench gas. Thus, in expressions such as Eq. (6) and Eq. (7) we use the instantaneous one photon detuning $D$,

$$\Omega_+ \rho_{01} = \text{const} - \frac{i|\Omega_+|^2}{4(\gamma_2 - iD)} d - \frac{i\Omega_+ \Omega_+^*}{2(\gamma_2 - iD)} \rho_{21}$$

(11)
To simplify this discussion and compare with experiment, we take the light fields to be real and equal, $\Omega_+ = \Omega_- = \Omega$ and find,

$$
\delta I_+ = \frac{\Omega^2 \gamma_2}{4(\gamma_2^2 + D^2)} \left[ 1 + \frac{\Omega^2}{8(\gamma_2^2 + D_0^2)(\Gamma^2 + \Delta^2)} \times \left( 4(\gamma_2 \Gamma' + D \Delta) - \frac{\Omega^2 \gamma_2 D_0 \Delta}{(\gamma_2^2 + D_0^2) \Gamma_1} \right) \right]
$$

$$
\delta I_- = \frac{\Omega^2 \gamma_2}{4(\gamma_2^2 + D^2)} \left[ 1 + \frac{\Omega^2}{8(\gamma_2^2 + D_0^2)(\Gamma^2 + \Delta^2)} \times \left( 4(\gamma_2 \Gamma' - D \Delta) + \frac{\Omega^2 \gamma_2 D_0 \Delta}{(\gamma_2^2 + D_0^2) \Gamma_1} \right) \right]
$$

The intensity noise in each polarization channel, Eq. (13) and Eq. (14), is dominated by the laser’s instantaneous one-photon detuned light scattering off the slowly evolving ground state coherences/population differences. With this we can now assemble the products in the correlation statistic, Eq. (1), by ensemble averaging over $D$.

Directly from Eq. (13) and Eq. (14) one can see that ensemble averaging reproduces the key features of the experimental data as seen in the representative $g^2(0)$ EIT intensity correlation curve in Fig. 2. Anticorrelation arises from the difference in signs in the last two terms of these two equations. These terms vanish as $\Delta \to 0$, which leads to the central correlation peak at zero two-photon detuning. At large $\Delta$ these anticorrelative terms vanish via their $\Omega^2 \gamma_2 \Gamma_1$ prefactor, and thus the naive power broadened EIT width $\Gamma'$ sets the scale of the outer envelope. This concludes our introduction to the simplest model that quantitively characterizes EIT noise spectroscopy.

V. COMPARISON: UNDERSTANDING THE NOISE CORRELATION POWER BROADENING REGIME

Experiment indicates that at high optical powers the central EIT intensity correlation peak broadens. This result is a challenge to the naive theory described above as numerical evaluation does not include a priori predictions of this behavior.
There are, however, physical causes of this new power broadening regime. As an example, a relationship between the laser’s RAM and power broadening of the correlation peak has been established experimentally and modelled theoretically in [14]. RAM within the power-broadened EIT bandwidth is quasi-adiabatically followed by the ground state coherences. Thus subsequent one-photon scattering (a fast process) off this noisy ground state coherence (a slow process) leads to a sampling over different, complex ground state coherences. This sampling creates an effective reduction in the norm of the coherence $\rho_{12}$ relative to average value of the norm computed via substituting $D_0$ into Eq. (4), Eq. (5) and Eq. (8). Referring back to the EIT correlation peak power-broadening experiments described in Section III and shown in Fig. 4 we can understand these results in terms of this RAM-assisted coherence reduction, as our laser’s fractional RAM was roughly 0.2%.

However, other physical process not reducible to parameters in the steady state model of the system that lead to depolarization of the ground state coherences will also contribute to power broadening of the correlation resonance [26]. One can see this directly in Eq. (9) and Eq. (10) and the equations that follow from them; as $\Delta \to 0$, $\rho_{12}$ becomes real and large while $\Delta$ becomes purely imaginary and small. Therefore any process that tends to rescale the coherences towards zero makes reaching anticorrelation in $g^2(0)$ harder to achieve. The main conclusion we use from the earlier studies cited above is that the power broadening slope is proportional to additional ground state coherence depolarization. Other physical processes leading to this depolarization that we have studied include imperfect bandwidth-limited averaging in the two-photon detuning map across the cell and coherence diffusion in the vapor cell.

Let the longitudinal distribution of two-photon detunings across the cell have a width $w$ and an average of $\Delta_0$. Atoms at various locations across the cell are described individually by Eq. (8) with a range of $\Delta$’s, all in the same rotating frame. The algebraic sum of these equations leads to,

$$\partial_t \bar{\rho}_{12} = -(\Gamma + i\bar{\Delta})\bar{\rho}_{12} + \frac{\bar{I}}{D_0} + d\rho_{12} d\Delta + \ldots \quad (15)$$

where $\bar{\cdot}$ represents the static ensemble average over the two-photon detuning map of the beam/cell, $\bar{I}/D$ is a catch-all for the intensity and detuning terms that are of the type already in Eq. (8) and the $d\rho_{12} d\Delta$ term is the convolution of the two-photon detuning map with the deviations of the ground state coherence across the beam/cell. The two-photon detuning map is characterized by $w$ and $\Delta_0$ so that to leading order the convolution yields
a term proportional to $\bar{\rho}_{12}$. As long as the variance is small compared to the $\Gamma$, when included via Eq. (15) in the steady state evaluation of Eq. (8), these deviations in the two-photon detuning across the cell effectively reduce the net ground state coherence $\bar{\rho}_{12}$ by an amount proportional to the ratio of the variance to the ground state decoherence rate, $w/\Gamma$. In simulation this reduction in the ground state coherence leads to a noise correlation resonance power-broadening slope also proportional to $w/\Gamma$.

This theory result is consistent with the observed change in the power broadening slope with applied $B$-field gradient as described in Section III and shown in Fig. 5. That is, a larger $B$-field gradient generates a two-photon detuning map with a larger variance, $w$, leading to a proportionally faster broadening rate. The results shown in Fig. 4(b) are also broadly consistent with this theory result, as the higher pressure buffer gas cell yields smaller $\Gamma$, largely determined by the diffusion rate out of the beam, $\Gamma_{df}$. However, as discussed in Section III another prominent effect of buffer gas pressure is collisional broadening of the excited state, and our numerical simulations show that the resulting difference in the effective Rabi frequencies accounts for the broadening rate discrepancy between the two buffer gas cells, in experiments where no substantial $B$-field gradient was applied.

To summarize, our studies indicate that all other things being equal (including the variance in the detuning map, other major causes of noise correlation resonance power broadening, and effective Rabi frequencies), the cell with the smaller intrinsic linewidth $\Gamma$ is expected to have a noise correlation resonance that power broadens more quickly at large optical powers.

VI. CONCLUSIONS AND ACKNOWLEDGEMENTS

In conclusion, we have demonstrated that power broadening of the EIT noise correlation peak is a consequence of processes that reduce the ground state coherences relative to that expected in the steady state for a single atom. Externally applied noise in the EIT bandwidth (like RAM) and imperfect averaging over two-photon detunings (for example, caused by a spatial inhomogeneity) are examples of processes that contribute to noise correlation power broadening. Theory and experiment both indicate that the power broadening slope is proportional to the inhomogeneous broadening associated with a variegated two-photon detuning environment and in that case is inversely proportional to the intrinsic two-photon
The narrow coherence resonance arising from EIT-derived intensity correlations using a low-cost non-narrow linewidth laser has broad applicability due to the simple nature of coupling intrinsic laser phase noise to ground state coherence evolution. Understanding the contributions to power broadening of the coherence peak is essential for broadening mitigation, which is desirable in technical applications. Further, this noise spectroscopy technique could also be used in other coherent media including EIT in rare-earth doped crystals [27], nitrogen-vacancy color centers in diamond [28], and quantum wells [29].

We are grateful to D. Phillips, R. Walsworth, A. Leandhardt, T. Walker, J. R. Brandenberger and Y. Xiao for useful discussions and equipment use. Additionally, SO would like to thank M. V. Camp, A. Stevenson, and H. Strehlow for their early contributions to this work.

This research was supported by an award from Research Corporation for Science Advancement (SO). MC acknowledges support under NSF grant number ECCS-1360725 and from the Science and Technology Center for Layered Polymeric Systems under grant number DMR 0423914.

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