Slow-light probe of Fermi pairing through an atom-molecule dark state

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We consider the two-color photoassociation of a quantum degenerate atomic gas into ground-state diatomic molecules via a molecular dark state. This process can be described in terms of a Λ level scheme that is formally analogous to the situation in electromagnetically induced transparency in atomic systems and therefore can result in slow-light propagation. We show that the group velocity of the light field depends explicitly on whether the atoms are bosons or fermions, as well as on the existence or absence of a pairing gap in the case of fermions, so that the measurement of the group velocity realizes a nondestructive diagnosis of the atomic state and the pairing gap.

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I. INTRODUCTION

Degenerate atomic Fermi gases have attracted much interest in recent years, well past the confines of traditional atomic, molecular, and optical (AMO) physics [1]. The existence of correlated Fermi pairs results in a number of effects that can be explored particularly well in these systems, due in particular to the control of two-body interactions provided by Feshbach resonances. These include detailed studies of the crossover from Bardeen-Cooper-Schrieffer (BCS) superfluidity to Bose-Einstein condensation (BEC) [1], of crystalline and supersolid phases [2], as well as spin-charge separation or spin drag [3], to mention a few examples. However, in the absence of any obvious change of density profile, the detection of Fermi pairing is challenging, in sharp contrast to the familiar BEC transition of bosons. A long-standing goal remains therefore to develop methods to efficiently detect the pairing signature of fermionic systems and other related exotic phases. Approaches toward this goal have focused on the measurement of atomic density-density correlations via the resonant or nonresonant optical response of the fermionic atoms [4], including methods of radio-frequency spectroscopy [5], photoemission spectroscopy [6], and Raman spectroscopy [7]. Alternative methods, such as scanning tunneling microscopy [8] or acoustic attenuation [9], are also being actively pursued.

In parallel to these developments, rapid experimental advances have resulted in the coherent formation of ultracold molecules from Bose or Fermi atoms [10]. The stable formation of diatomic molecules from laser-cooled alkali-metal atoms has been achieved by using magnetic Feshbach resonances and optical photoassociation (PA) techniques. By applying an all-optical PA method, molecules associated with ultracold atoms can be successfully transferred into their rovibrational ground state [11].

A key component of the two-color PA method is the existence of an atom-molecule dark state, as first demonstrated by Winkler *et al.* [12]. The underlying quantum interference and slow-light propagation were also observed for ultracold sodium atoms by Turner *et al.* [13], hinting at the possibility of studying the quantum control of light through cold reactions [10–14], quantum state transfer from light to molecules [14,15], as well as high-precision diagnostics of Fermi gases via PA spectroscopy [7].

In this paper we show that the slow-light propagation associated with the existence of such a dark state also provides a relatively simple nondestructive probe of Fermi pairing, without the need for additional excitations (atom-to-atom, atom-ion-to-molecule, or molecule-to-molecule) or for laser imaging of the populations of transferred particles. This proposed method finds additional motivation in previous work [16] that showed that the statistical properties of the molecular field formed from ultracold atoms depends strongly on the statistical properties of these atoms. In particular, it was found that for short times, the number of molecules created scales as the square N^2 of the number of atoms in the case of an atomic Bose-Einstein condensate, but as N for a normal Fermi gas at zero temperature, a manifestation of the independence of all atomic pairs in that case. For a paired Fermi gas, the situation is intermediate between these two extremes: the molecules are formed at a higher rate than for a normal Fermi gas, and the maximum number of molecules is larger, approaching the BEC situation for strong pairing.

The main result of the present analysis is that a related situation does occur when considering the dark-state propagation of a photoassociating light field: in contrast to the case where photoassociation originates from a condensate of bosonic atoms, and where the inverse group velocity v_g^{-1} of the light field is known to scale as N^2, we find that for a normal Fermi gas at $T = 0$ it scales as N. A paired Fermi system represents an intermediate situation, as was the case in Ref. [16]. It follows that the group velocity is a direct measure of the pairing gap Δ. This simple all-optical method is also expected to prove useful in probing, e.g., polaron-to-molecule transitions and atom-molecule vortex states [17] by photoassociating a spin-imbalanced or a rotating Fermi gas. We remark that this proposal involves the use of tunable atom-molecule interactions and as such is fundamentally different from approaches based on single-atom excitations [4,18].

The paper is organized as follows. Section II describes our model and calculates the slow-light group velocity of a quantized optical field that propagates in a normal Fermi gas and helps photoassociating atoms into molecules via a dark-state intermediate level. Section III evaluates the effect of a

Fermi pairing gap on that velocity and shows that it depends strongly on the magnitude of the gap. Finally Sec. IV is a conclusion and outlook.

II. NORMAL FERMI GAS

We first consider the two-color photoassociation of a homogeneous, normal degenerate Fermi gas with no pairing. The entrance channel atoms, the intermediate state $|m\rangle$, and the closed-channel bosonic molecules are characterized by the annihilation operators $\hat{c}_{\mathbf{k}\sigma}$, $\hat{m}_{\mathbf{k}+\mathbf{k}'}$ and $\hat{a}$, respectively, where $\mathbf{k}$ and $\mathbf{k}'$ are wave numbers and σ labels the fermionic spin. We assume that the PA between atomic pairs and excited molecules in state $|m\rangle$ is driven by an optical field that is treated quantum mechanically at that point, and the field that drives the molecules to their ground state $|g\rangle$ is classical, with Rabi frequency $\Omega(t)$ (see Fig. 1).

At the simplest level, the Hamiltonian of this system can be expressed as ($\hbar = 1$)

$$\hat{H} = \sum_{\mathbf{k},\sigma} \frac{\epsilon_{\mathbf{k}}}{2} \hat{c}^{\dagger}_{\mathbf{k}\sigma} \hat{c}_{\mathbf{k}\sigma} + g \sum_{\mathbf{k},\mathbf{k}'} (\hat{\mathcal{E}} \hat{m}^{\dagger}_{\mathbf{k}+\mathbf{k}'} \hat{c}_{\mathbf{k}\uparrow} \hat{c}_{\mathbf{k}'\downarrow} + \text{H.c.}) + \sum_{\mathbf{k},\mathbf{k}'} [\delta \hat{m}^{\dagger}_{\mathbf{k}+\mathbf{k}'} \hat{m}_{\mathbf{k}+\mathbf{k}'} + \Omega(\hat{a} \hat{m}^{\dagger}_{\mathbf{k}+\mathbf{k}'} + \text{H.c.})], \quad (1)$$

where g is the atom-molecule coupling constant, δ is the detuning between the frequency of the quantized photoassociation field and the frequency difference between the atomic fermions and the molecular state $|m\rangle$ (we neglect the dispersion in fermionic energies $\epsilon_{\mathbf{k}}$ for simplicity), and $\Omega(t)$ is the Rabi frequency of the classical field, taken to be real without lack of generality. The s-wave collisions between fermionic atoms, between molecules, and between atoms and molecules are ignored for a dilute gas. We note that a quantized description of the photoassociating field is not required in the calculation of its group velocity, but as further discussed later, it demonstrates the potential to achieve control of the statistics of the closed-channel molecules.

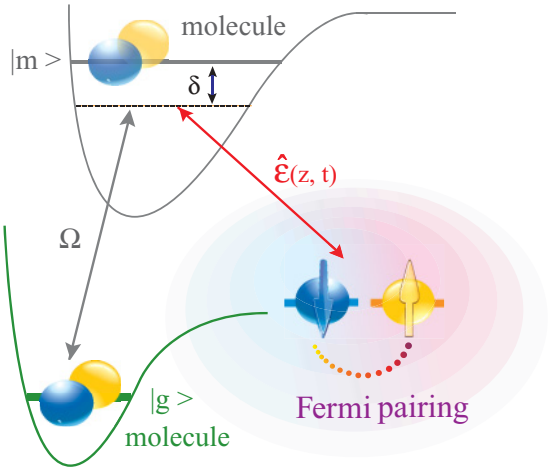

FIG. 1. (Color online) Schematic of two-color PA in an ultracold degenerate Fermi gas with or without Cooper pairing. The Λ configuration illustrated here highlights the close formal analogy to the familiar situation of atomic EIT [18], but any narrow intermediate molecular level will work as well in the weak probe limit.

We restrict ourselves for simplicity to the association of atom pairs with opposite momenta ($\mathbf{k} = -\mathbf{k}'$) and opposite spin, in which case the intermediate molecules can be also described in terms a single-mode bosonic field when concentrating on short-time dynamics; see, e.g., Refs. [16,19]. With these simplifying assumptions, this system is formally analogous to the situation of electromagnetically induced transparency (EIT) in atomic Λ systems and can be expected to result in slow-light propagation.

We describe the propagation of the light fields in the atomic ensemble in an effective one-dimensional theory, following the approach familiar from nonlinear optics and fiber optics. We emphasize that this does not imply that we consider a one-dimensional atomic sample. Rather, the effective one-dimensional propagation of the light field results from an integral over the transverse dimensions of the sample; see, e.g., Ref. [20]. This approach is appropriate provided the Rayleigh length of the light fields is large compared to the length of the atomic sample, and that focusing and defocusing effects can be neglected.

The quantized optical field $\hat{E}(z,t)$ of carrier frequency ν is then given by

$$\hat{E}(z,t) = \sqrt{\frac{\hbar\nu}{2\epsilon_0 L}} \hat{\mathcal{E}}(z,t) \exp[i\nu(z/c - t)],$$

where L is the quantization length. It satisfies the commutation relation

$$[\hat{E}(z,t), \hat{E}^{\dagger}(z',t)] = \frac{\nu}{\epsilon_0} \delta(z - z').$$

Within the slowly-varying-amplitude approximation, the propagation equation of the field envelope $\hat{\mathcal{E}}(z,t)$ is given by

$$\left(\frac{\partial}{\partial t} + c\frac{\partial}{\partial z}\right) \hat{\mathcal{E}}(z,t) = igL \sum_{\mathbf{k}} \hat{c}^{\dagger}_{-\mathbf{k}\downarrow}(z,t) \hat{c}^{\dagger}_{\mathbf{k}\uparrow}(z,t) \hat{m}(z,t). \quad (2)$$

In the following we consider the regime of weak excitations, where the atomic population remains essentially undepleted. The initial state of the atom-molecule system is taken as

$$|\psi(0)\rangle = |F\rangle \otimes |0\rangle_m \otimes |0\rangle_a,$$

where $|0\rangle_m$ and $|0\rangle_g$ denote the vacuum state for the molecules,

$$|F\rangle = \prod_k \hat{c}^{\dagger}_{-\mathbf{k}\downarrow} \hat{c}^{\dagger}_{\mathbf{k}\uparrow} |0\rangle,$$

and the product is taken up to the Fermi surface, a step appropriate for temperatures much below the Fermi temperature [16]. Introducing the pseudo-spin operators

$$\begin{aligned} \hat{s}^{+}_{\mathbf{k}} &= (\hat{s}^{-}_{\mathbf{k}})^{\dagger} = \hat{c}^{\dagger}_{-\mathbf{k}\downarrow} \hat{c}^{\dagger}_{\mathbf{k}\uparrow}, \\ \hat{s}^{z}_{\mathbf{k}} &= \tfrac{1}{2}(\hat{c}^{\dagger}_{\mathbf{k}\uparrow} \hat{c}_{\mathbf{k}\uparrow} + \hat{c}^{\dagger}_{-\mathbf{k}\downarrow} \hat{c}_{-\mathbf{k}\downarrow} - 1), \end{aligned} \quad (3)$$

which satisfy the commutation relations

$$[\hat{s}^{+}_{\mathbf{k}}, \hat{s}^{-}_{\mathbf{k}'}] = 2\delta_{\mathbf{k}\mathbf{k}'} \hat{s}^{z}_{\mathbf{k}}, \quad [\hat{s}^{z}_{\mathbf{k}}, \hat{s}^{\pm}_{\mathbf{k}'}] = \pm\delta_{\mathbf{k}\mathbf{k}'} \hat{s}^{\pm}_{\mathbf{k}}, \quad (4)$$

and the collective operators

$$\begin{aligned} \hat{S}_{\pm} &= \sum_{\mathbf{k}} \hat{s}^{\pm}_{\mathbf{k}}, \quad \hat{S}_z = \sum_{\mathbf{k}} \hat{s}^{z}_{\mathbf{k}} = \frac{N}{2} - \hat{a}^{\dagger}\hat{a} - \hat{m}^{\dagger}\hat{m}, \\ \hat{\mathbf{S}}^2 &= \hat{S}_+ \hat{S}_- + \hat{S}_z(\hat{S}_z - 1), \end{aligned} \quad (5)$$

with the conserved total number of atomic pairs and molecules

$$N = \sum_k (\hat{c}^\dagger_{\mathbf{k}\uparrow}\hat{c}_{\mathbf{k}\uparrow} + \hat{c}^\dagger_{-\mathbf{k}\downarrow}\hat{c}_{-\mathbf{k}\downarrow})/2 + (\hat{a}^\dagger\hat{a} + \hat{m}^\dagger\hat{m})$$
$$= (\hat{S}_z + N/2) + (\hat{a}^\dagger\hat{a} + \hat{m}^\dagger\hat{m}), \quad (6)$$

yields for the Hamiltonian $\hat{H}_\mathcal{N}$ the simplified form

$$\hat{H} = \sum_{\mathbf{k}} \epsilon_{\mathbf{k}}\hat{s}^z_{\mathbf{k}} + \delta\hat{m}^\dagger\hat{m} + (g\hat{\mathcal{E}}\hat{m}^\dagger\hat{S}_- + \Omega\hat{m}^\dagger\hat{a} + \text{H.c.}). \quad (7)$$

The resulting Heisenberg equations of motion are, by approximating all $\epsilon_{\mathbf{k}}$'s by the Fermi energy ϵ_F,

$$i\frac{d\hat{S}_z}{dt} = g\hat{\mathcal{E}}^\dagger\hat{m}\hat{S}_+ - g\hat{\mathcal{E}}\hat{m}^\dagger\hat{S}_-, \quad i\frac{d\hat{S}_-}{dt} = \epsilon_F\hat{S}_- - 2g\hat{\mathcal{E}}^\dagger\hat{m}\hat{S}_z,$$
$$i\frac{d\hat{S}_+}{dt} = -\epsilon_F\hat{S}_- + 2g\hat{\mathcal{E}}\hat{m}^\dagger\hat{S}_z, \quad i\frac{d\hat{m}}{dt} = g\hat{\mathcal{E}}\hat{S}_- + \delta\hat{m} + \Omega\hat{a},$$
$$i\frac{d\hat{a}}{dt} = \Omega\hat{m}, \quad i\frac{d\hat{\mathcal{E}}}{dt} = g\hat{m}\hat{S}_-. \quad (8)$$

In the following we consider the resonant situation $\delta = 0$. In the limit of weak excitations considered here we can also set $d\hat{m}/dt \to 0$. We then have, in the lowest nonvanishing order in the excited molecular state [14,21],

$$\hat{a} = -(g/\Omega)\hat{\mathcal{E}}\hat{S}_-, \quad \hat{m} = -i(g/\Omega)\hat{S}_-\frac{\partial}{\partial t}\left(\frac{\hat{\mathcal{E}}}{\Omega}\right). \quad (9)$$

The propagation of the field $\hat{\mathcal{E}}(z,t)$ is then governed by the equation

$$\left(\frac{\partial}{\partial t} + c\frac{\partial}{\partial z}\right)\hat{\mathcal{E}}(z,t) = -\frac{g^2LN}{\Omega}\frac{\partial}{\partial t}\left(\frac{\hat{\mathcal{E}}}{\Omega}\right), \quad (10)$$

where we have used

$$\hat{\mathbf{S}}^2|F\rangle = S(S+1)|F\rangle = \frac{N}{2}\left(\frac{N}{2}+1\right)|F\rangle, \quad (11)$$

and the fact that in the weak excitation approximation we have

$$\langle\hat{S}_+\hat{S}_-\rangle = \left(-n_a^2 + n_aN - n_a\right) + N \sim N. \quad (12)$$

Equation (10) can be recast as

$$\left(\frac{\partial}{\partial t} + \frac{c}{1+\beta_f}\frac{\partial}{\partial z}\right)\hat{\mathcal{E}}(z,t) = \frac{\beta_f}{1+\beta_f}\left(\frac{1}{\Omega}\frac{\partial\Omega}{\partial t}\right)\hat{\mathcal{E}}. \quad (13)$$

where

$$\beta_f \equiv \frac{g^2LN}{\Omega^2}. \quad (14)$$

That is, the group velocity of the field $\hat{\mathcal{E}}(z,t)$ is

$$v_g = \frac{c}{1+\beta_f} = c\cos^2\theta, \quad (15)$$

with

$$\theta = \tan^{-1}(g\sqrt{LN}/\Omega). \quad (16)$$

As mentioned in the Introduction, the scaling of β_f with N should be contrasted with the situation for a pure condensate of bosonic atoms, in which case it is predicted that [14]

$$\beta_f \to \beta_b = \frac{g^2LN^2}{\Omega^2} = N\beta_f. \quad (17)$$

As was the case in the analysis of molecule formation of Ref. [16], this difference is due to the fact that for a Bose-Einstein condensate the photoassociation is a collective atomic effect, while in a normal Fermi gas the atom pairs act independently from each other.

As already indicated, the form of v_g is independent of whether the field $\hat{\mathcal{E}}(z,t)$ is treated classically or quantum mechanically (see the related experiment of Ref. [13]). The quantized description used here is primarily to facilitate a direct comparison with the bosonic atom-molecule system of Ref. [14]. Note, however, that Eq. (9) shows that the statistical properties of the closed-channel molecules are determined by the states of both the optical field and the Fermi atoms, hinting at the possibility of quantum control of the closed-channel molecules, e.g., by applying a squeezed PA field [14].

The next section expands these considerations to the case of a paired Fermi gas, which is then expected to represent an intermediate situation between these two extremes. We show that this is indeed the case, and as a result, measuring the group velocity of the photoassociating field provides a direct measure of the pairing gap.

III. PAIRING AND GROUP VELOCITY

To account for the impact of Cooper pairing on the group velocity v_g we include attractive pairing interactions into Eq. (1) in the usual fashion via the Hamiltonian [16,19]

$$\hat{H}_{\text{BCS}} = \hat{H} - U\sum_{k,k'}\hat{s}^+_k\hat{s}^-_{k'}. \quad (18)$$

The BCS ground state is found as usual by minimizing $\langle\hat{H}_{\text{BCS}} - \mu\hat{N}\rangle$, where μ is the chemical potential, using the ansatz

$$|\text{BCS}\rangle = \prod_k(u_{\mathbf{k}} + v_{\mathbf{k}}\hat{s}^+_{\mathbf{k}})|0\rangle, \quad (19)$$

with the result

$$\begin{pmatrix} u^2_{\mathbf{k}} \\ v^2_{\mathbf{k}} \end{pmatrix} = \frac{1}{2}\left(1 \mp \frac{\xi_{\mathbf{k}}}{\sqrt{\xi^2_{\mathbf{k}} + |\Delta|^2}}\right). \quad (20)$$

Here $\eta_{\mathbf{k}} = \sqrt{\xi^2_{\mathbf{k}} + |\Delta|^2}$ is the mean-field quasiparticle energy, $\xi_{\mathbf{k}} = \epsilon_{\mathbf{k}} - \mu$ is the kinetic energy of the atoms measured from the Fermi surface, and

$$\Delta = U\sum_{\mathbf{k}} u_{\mathbf{k}}v_{\mathbf{k}} = \frac{U}{2}\sum_{\mathbf{k}}\frac{\Delta}{\sqrt{\xi^2_{\mathbf{k}} + |\Delta|^2}} \quad (21)$$

is the gap parameter.

The interaction Hamiltonian (18) does not modify the equations of motion for the operators $\hat{m}$, $\hat{a}$, and $\hat{\mathcal{E}}$. In the present context, its main effect in the weak excitation limit is to replace $\langle\hat{S}_+\hat{S}_-\rangle$ by

$$\langle\hat{S}_+\hat{S}_-\rangle = \sum_{\mathbf{k}} v^2_{\mathbf{k}} + \sum_{\mathbf{k}\neq\mathbf{k}'} u_{\mathbf{k}}v_{\mathbf{k}}u_{\mathbf{k}'}v_{\mathbf{k}'} \simeq N + \left(\frac{\Delta}{U}\right)^2. \quad (22)$$

Within the weak-coupling limit of BCS theory, $\epsilon_{\mathbf{k}}$ and $\xi_{\mathbf{k}}$ are approximately independent of the wave vector $\mathbf{k}$, $\epsilon_{\mathbf{k}} \to \epsilon_F$

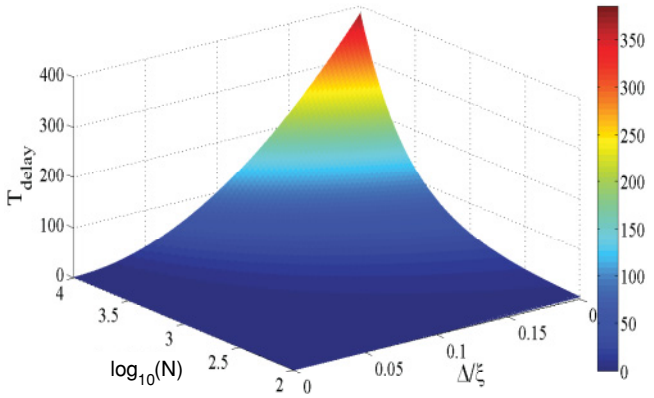

FIG. 2. (Color online) Dimensionless relative time delay T_d (scaled by L/v_g) as a function of the dimensionless pairing gap Δ/ξ and the total number of atomic pairs and molecules N.

and $\xi_{\mathbf{k}} \to \xi$, where ϵ_F is the Fermi energy [22]. In that case the group velocity becomes

$$v_{g,\Delta} = \frac{c}{1+\beta_\Delta}, \tag{23}$$

where

$$\beta_f \to \beta_\Delta = \beta_f \left(1 + \frac{N\Delta^2}{4\xi^2 + 4\Delta^2}\right), \tag{24}$$

indicating that it now depends on both N *and* the pairing gap Δ.

This is illustrated in Fig. 2, which shows the time delay

$$T_d = \frac{L}{v_{g,\Delta}} - \frac{L}{v_g} = \frac{L\beta_\Delta}{c} \tag{25}$$

experienced by a short photoassociating light pulse as a function of N and the pairing gap Δ, relative to the delay in the absence of gap. For large values of Δ, we have $v_{g,\Delta} \sim N^{-2}$, approaching the case of a bosonic atom-molecule dark-state medium [14], as would be intuitively expected. The gap-dependent enhancement factor that is determined precisely by the ratio of the molecule population $N_a(\Delta)$ and N_a in the presence or absence of a pairing gap,

$$\zeta = 1 + \frac{N\Delta^2}{4(\xi^2+\Delta^2)} = \frac{N_{a,\Delta}}{N_a}, \tag{26}$$

see Fig. 3. That is, the variation in group velocity originates directly from the PA-induced atom-molecule superpositions in the Λ level scheme of Fig. 1.

Note that it is the relative number of molecules (i.e. the ratio of the cases with and without BCS pairing) that is plotted in Fig. 3. That ratio approaches unit for very small pairing gaps, as should be expected. The absolute number of molecules, however, is proportional to $(g/\Omega)^2 \sim 10^{-4}$ for typical values $g \sim 10$ kHz, $\Omega \sim 1$ MHz, and thus remains quite small in comparison with the total number of atoms. This is again similar to the situation in purely atomic systems, where the weak excitation approximation was shown to be valid in the regime of weak probe fields [18].

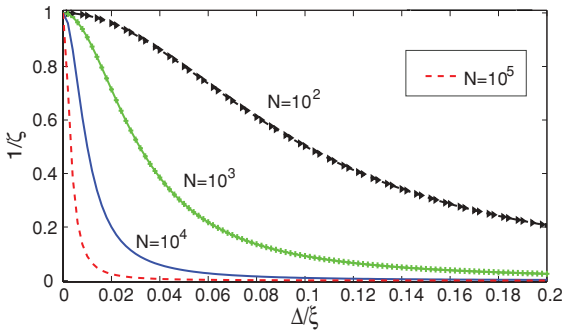

FIG. 3. (Color online) Relative molecule population $\zeta^{-1} = N_a/N_a(\Delta)$ as a function of the dimensionless pairing gap Δ/ξ for several values of N. The outer-most curve corresponds to $N = 10^2$ and the inner-most to $N = 10^5$.

IV. CONCLUSION

In conclusion, we have shown that the two-color photoassociation of fermionic atoms into bosonic molecules via a dark-state transition results in a group velocity of the photoassociating field that can be slowed significantly, in complete analogy with the situation of EIT in Λ three-level atomic systems. The scaling of the group velocity v_g with particle number depends not only on whether the atoms are bosonic or fermionic, with an associated N^2 versus N dependence, but also on the possible pairing of the fermionic atoms resulting from attractive two-body interactions. As such, the propagation delay of the photoassociating light pulse $\hat{\mathcal{E}}(z,t)$ provides a direct measurement of the pairing gap Δ. This nondestructive *in situ* diagnostic technique supports and extends the idea of using Raman spectroscopy [7] to extract the pairing parameters, but differs from proposals based solely on the use of atomic transitions [18]. We note that the different features in slow-light propagation for both bosonic and fermionic atoms were also studied previously for purely atomic transitions (without any photoassociation) [23].

To estimate the pairing-induced optical time delay of the propagating pulse, we consider the typical values $g \sim 10$ kHz, $\Omega \sim 1$ MHz, $N = 10^6$, $L = 1$ mm, and $\gamma_m \sim 16$ MHz, $\gamma_a \sim 600$ Hz [7]. These values give for the bosonic sample a group velocity of $v_g \sim 3$ km s^{-1}, that is, a significant slowing down of the light pulse. For a normal Fermi gas, the significantly less favorable scaling of v_g with N instead of N^2 gives $v_g \sim 0.9c$ for the same parameters. The extremely small change in propagation time compared to the vacuum case will be challenging to measure. Even for a relatively long sample of 1 mm, the difference in propagation times will be of the order of 0.3×10^{-12} s, a delay that might, however, be observable by an interferometric technique. For paired fermionic atoms, finally, we find $v_g \sim 300$ km s^{-1} for $\Delta/\xi = 0.2$, and $v_g \sim 15$ km s^{-1} for $\Delta/\xi = 2$, a change of three to four orders of magnitude compared to the case of a normal Fermi gas. Such a change should be easily observable. Note that shorter samples lead to a reduction in delay time T_d that scales as L^2, as readily seen from Eqs. (14) and (25), rapidly leading to significant reductions in time delay and increasingly challenging experiments.

Our discussion ignores the decay of molecular states. However, it can be readily shown that after including these

decay terms, the group velocity of the signal is still in the form of Eq. (15), but with the substitution $\Omega \to \sqrt{\Omega^2 + \gamma_m \gamma_a}$ [14]. In practice, the PA pulse duration τ should satisfy $\tau \ll \gamma_a^{-1} \sim 1.67$ ms, a condition that can be fulfilled in current experiments [7,24–26].

Future work will improve the sample description by incorporating its transverse effects in the propagation of light fields, finite temperature effects, and a more detailed description of the two-body physics. In this context it will also be interesting to consider cavity-induced transparency with a degenerate Fermi gas [27]. A significantly more challenging problem will involve the situation of strong pair fluctuations at the BEC-BCS crossover region [28,29]. Finally, we note that the use of nonclassical associating light fields will also allow one to consider the correlations of the transmitted field and/or a possible molecule-photon entanglement as probes of Fermi pairing or possibly of other exotic phases.

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