Feshbach resonance-induced Fano interference in photoassociation

Bimalendu Deb and G S Agarwal

1 Department of Materials Science, and Raman Center for Atomic, Molecular and Optical Sciences, Indian Association for the Cultivation of Science (IACS), Jadavpur, Kolkata 700 032, India
2 Department of Physics, Oklahoma State University, StillWater, OK 74078, USA

E-mail: msbd@iacs.res.in

Received 14 August 2009, in final form 23 September 2009
Published 27 October 2009
Online at stacks.iop.org/JPhysB/42/215203

Abstract

We consider photoassociation from a state of two free atoms when the continuum state is close to a magnetic field-induced Feshbach resonance and analyse Fano interference in photoassociation. We show that the minimum in photoassociation profiles characterized by the Fano asymmetry parameter $q$ is independent of laser intensity, while the maximum explicitly depends on laser intensity. We further discuss the possibility of the nonlinear Fano effect in photoassociation near a Feshbach resonance.

(Some figures in this article are in colour only in the electronic version)

1. Introduction

In recent times, quantum interferences have occupied a prominent place in physics and these occur rather ubiquitously. Many well-known examples of these include Fano interferences [1–3], electromagnetically induced transparency (EIT) [4], vacuum-induced interferences in spontaneous emission [5]. The quantum interferences have resulted in a large number of applications in coherent control of the optical properties, control of spontaneous emission [6] and slow light [7, 8]. Quantum interference has been experimentally demonstrated in the coherent formation of molecules [9] and Autler–Townes splitting [10, 11] in two-photon PA. Theoretical formulation of PA within the framework of Fano’s theory has been developed in [12] and [13]. Recent experimental [14–16] and theoretical [17–20] studies on photoassociation (PA) near a magnetic field Feshbach resonance (MFR) [21] have generated a lot of interest in Fano interference with ultracold atoms. In a remarkable experiment, Junker et al [14] have demonstrated asymmetric spectral line shape and saturation in PA due to a tunable MFR. Asymmetric line shape is a hallmark of the Fano effect and the experimental results of [14] can be attributed to the Fano interference.

Here we demonstrate quantum interference in the context of photoassociation (PA) [22] under the condition when a Feshbach resonance is also involved in photoassociation. We show a Fano-like interference minimum in the photoassociation spectrum. In analogy to the well-known Fano $q$-parameter we can introduce a parameter which governs the existence of this minimum. Although the minimum is independent of laser intensity, the maximum is shown to depend explicitly on laser intensity. From our calculations we extract line shapes which are in broad agreement with the experimental results of Junker et al [14]. Our formula for photoassociation is expressed in terms of parameters each of which has a clear physical meaning and is measurable. We derive the probability of PA excitation for arbitrary intensities of the laser field and thus we also discuss the nonlinear Fano effect. The current work has some features in common with the recent paper of Kuznetsova et al [20] though these authors address a different problem which is the population transfer using two laser beams. Our emphasis is on quantum interferences in PA using a single laser beam.

The paper is organized in the following way. In section 2, we consider a simple model of three-channel time-independent scattering in the presence of an optical and a magnetic field. By using Green’s functions, we present a compact analytical solution of the model. We then discuss selective results in section 3. The paper is concluded in section 4.

2. The model and its solution

To begin with, we model PA in the presence of a Feshbach resonance as a three-channel scattering problem. There are
two ground-state asymptotic hyperfine channels of which one is closed and the other one is open. The third channel corresponds to the photoassociated excited molecular configuration. The two ground-state channels are coupled via the hyperfine interaction. At a Feshbach resonance, the two atoms will form a quasibound state in the closed channel as schematically illustrated in figure 1. As the strength of the applied magnetic field is varied, this quasibound state can move across the collision energy. When a PA laser is applied to form an excited photoassociated molecule (PM), there arise two competing pathways of dipole transitions as shown by different colours in figure 1. One is the continuum–bound and the other one is bound–bound transition. We assume that the energy spacing of closed-channel quasibound states and rotational spacing of PM states are much larger than PA laser linewidth so that only one rotational level \( J \) of a particular vibrational state \( v \) of PM is coupled to a particular quasibound state by the PA laser. This assumption is particularly good for relatively deeply bound vibrational levels with outer turning point lying at a separation less than 30 a₀ (a₀ is the Bohr radius). Usually, these levels have rotational spacings exceeding the PA laser linewidth by orders of magnitude. At ultracold temperatures (<1 mK), all the partial waves except \( \ell = 0 \) (s-wave) of the relative motion of two ground-state atoms are negligible. This means that, in the molecular frame of reference of two ground-state ultracold atoms, \( J = 0 \) rotational state is the most significant. When the Feshbach resonance leads to the association of two ultracold ground-state atoms into a bound state embedded in the continuum or a Feshbach molecule, the bound state or the molecule should also be predominantly in the non-rotating \( (J = 0) \) state. Under such conditions, the rotational selection rule of optical transition tells us that only \( J = 1 \) rotational state of the photoassociated molecule will be significantly populated. In our model calculations, the consideration of only one vibrational level in the closed-channel bound state can be justified by having a system which has non-overlapping Feshbach resonances and selecting one such resonance with resonance width \( \Gamma_f \) much larger than the laser linewidth.

Let us write an energy eigenstate of the system of two atoms interacting simultaneously with a magnetic and a PA laser field in the form

\[
|\Psi_E\rangle = \Phi_f|g_2\rangle + \chi|g_1\rangle + \Phi_p|e\rangle
\]  

where \( E \) is an energy eigenvalue, \(|g_1(2)\rangle\) represents the internal electronic states of 1 (2) or open (closed) channel and \(|e\rangle\) denotes the electronic state of the excited molecule. \( \Phi_f \) and \( \Phi_p \) are the diatomic bound states. The continuum state has the form \( \chi = \int dE' b_E \Psi_{E'} \) where \( \Psi_{E'} \) is an energy-normalized scattering state of collision energy \( E' \) and \( b_E \) is the density of unperturbed continuum states. The state (1) is assumed to be energy-normalized. The Hamiltonian of the system can be written as \( H = H_{\text{kin}} + H_{\text{elec}} + H_{\text{hfs}} + H_B + H_L \) where \( H_{\text{hfs}} \) denotes a term corresponding to the total kinetic energy of the two atoms and \( H_{\text{elec}} \) is a term that depends only on electronic coordinates of the two atoms, \( H_B \) is the hyperfine interaction term. Here \( H_B \) represents the magnetic interaction in the atomic states, and \( H_L \) represents the laser interaction between atomic or molecular states. From the time-independent Schrödinger equation \( H \Psi_E = E \Psi_E \) under Born–Oppenheimer approximation, one obtains the following coupled equations:

\[
\begin{align*}
\left[ -\frac{\hbar^2}{2\mu} \frac{d^2}{dr^2} + B_J(r) \right] \Phi_f + \left( V_J(r) - \hbar \omega_L - E - i\hbar \frac{\chi}{2} \right) \Phi_P &= -\Lambda_1 \chi - \Lambda_2 \Phi_f \\
\left[ -\frac{\hbar^2}{2\mu} \frac{d^2}{dr^2} + V_J(r) - E \right] \Phi_J &= -\Lambda_2 \Phi_P - V \Phi_J
\end{align*}
\]

where \( \omega_L \) is the laser frequency, \( \Lambda_1 \) and \( \Lambda_2 \) are the laser-induced transition dipole matrix elements between \(|e\rangle\) and \(|g_1\rangle\) and between \(|e\rangle\) and \(|g_2\rangle\), respectively. Here \( V_J \) (\( J = 1, 2 \)) are the potentials including hyperfine and Zeeman terms, \( V_J \) is the excited state molecular potential and \( V \) stands for the hyperfine spin coupling between closed-channel bound state and continuum states. Here \( B_J(r) = \hbar^2 J(J + 1)/(2\mu r^2) \) is the rotational term of the excited state. Note that for the ground scattering and bound states we have considered only the zero rotational state. The zero of energy scale is taken to be the threshold \( E_{\text{th}} \) of the open channel 1 and the energies of the bound states are measured from this reference. For two homonuclear atoms, the asymptotic form of the potential \( V_J(r \to \infty) \sim \hbar \omega_A - C_3/r^3 \), where \( C_3 \) is the long-range coefficient of the dipole–dipole interaction between one ground-state S-atom and another excited state P-atom and \( \omega_A \) is the atomic frequency. These three coupled equations can be solved exactly by the use of real space Green’s function as described below.

It is convenient to write \( V_e = \hbar \omega_A + \bar{V}_{ex} \), where \( \bar{V}_{ex}(r \to \infty) \sim -C_3/r^3 \). Let \( \Phi_p \) denote the bound state solution of the potential \( \bar{V}_{ex} \) and \( E_p \) be the corresponding bound state (negative) energy. Green’s function for the homogeneous part with \( \Lambda_1 = \Lambda_2 = 0 \) (i.e. without laser couplings) of (2) can be
written as

\[
G_p(r, r') = -\frac{1}{\hbar \delta + E - E_p + i\hbar \gamma/2} \phi_p(r) \phi_p(r')
\]  

where \( \delta = \omega_L - \omega_A \). Using this function, we can write down the solution of equation (2) in the form \( \Phi_p = A_p \phi_p(r) \) where

\[
A_p = \frac{\int \text{d}r' \{ A_1(r') \chi(r') + A_2(r') \phi_p(r') \} \phi_p(r')}{\hbar \delta + E - E_p + i\hbar \gamma/2}. 
\]

Similarly, with the use of Green’s function for the homogeneous part of (3), we have \( \Phi_f(r) = A_m \phi_f(r) \) where

\[
A_m = \frac{\int \text{d}r' \{ V^*(r') \chi(r') + A_p \chi(r') \phi_p(r') \} \phi_f(r')}{E - E_f} 
\]

where \( \phi_f(r) \) is the wavefunction and \( E_f \) is the energy of bound state in the closed channel in the absence of laser field. Now, we can express \( A_p \) in terms of integrals involving the continuum state \( \chi \) and molecular bound states \( \phi_p \) and \( \phi_f \). Then substituting \( \Phi_p \) and \( \Phi_f \) expressed in terms of \( A_p \) and \( \chi \) in equation (4) and making use of the relation \( \chi = \int dE b_E \psi_E \), we obtain

\[
\hbar^2 \frac{d^2}{dr^2} \psi_E(r) + \left[V(r) - E\right] \psi_E(r) = -\Lambda^\ast_1(r) \tilde{A}_p \phi_p(r) 
\]

where

\[
\tilde{A}_p = \frac{A_p (E - E_f) + A_{pf} \tilde{V}_{fc}}{D(E - E_f) - |A_{pf}|^2}. 
\]

Here \( D = \hbar \delta + E - E_p + i\hbar \gamma/2 \), \( \tilde{V}_{fc} = \int dr \phi_f(r) V(r) \psi_E(r) \), \( \tilde{A}_{pc} = \int dr \phi_p(r) A_1(r) \psi_E(r) \) and \( \tilde{A}_{2, pf} = \int dr \phi_p(r) A_2(r) \phi_f(r) \). Equation (8) can now be solved by constructing Green’s function with the scattering solutions of the homogeneous part (i.e. for \( \Lambda_1 = V = 0 \)). This Green’s function can be written as

\[
K(r, r') = -\pi \left[ \psi_E^{0, \text{reg}}(r) \psi_E^{0, \text{irr}}(r') + i \psi_E^{0, \text{reg}}(r) \psi_E^{0, \text{irr}}(r') \right],
\]

\[
(r' > r)
\]

\[
K(r, r') = -\pi \left[ \psi_E^{0, \text{reg}}(r) \psi_E^{0, \text{irr}}(r') + i \psi_E^{0, \text{reg}}(r) \psi_E^{0, \text{irr}}(r') \right],
\]

\[
(r' < r)
\]

where the regular function \( \psi_E^{0, \text{reg}}(r) \) vanishes at \( r = 0 \) and the irregular solution \( \psi_E^{0, \text{irr}}(r) \) is defined by boundary only at \( r \to \infty \). These have the familiar asymptotic behaviour \( \psi_E^{0, \text{reg}}(r) \propto j_0 \cos \eta_0 - n_0 \sin \eta_0 \) and \( \psi_E^{0, \text{irr}}(r) \propto n_0 \cos \eta_0 + j_0 \sin \eta_0 \), where \( j_0 \) and \( n_0 \) are the spherical Bessel and Neumann functions for \( \ell = 0 \) and \( \eta_0 \) is the s-wave phase shift in the absence of laser and magnetic field couplings. Here \( E = \hbar^2 k^2/(2m) \) with \( m \) being the reduced mass of the two atoms. Next, we can express the solution of (8) in the following form:

\[
\psi_E = \exp(i\eta_0) \psi_E^{0, \text{reg}} + \frac{\int \text{d}r' K(r, r') \left[ \Lambda_1^\ast(r') \tilde{A}_p \phi_p(r') \right]}{E - E_f} + \frac{\tilde{V}_{fc} + \tilde{A}_p \Lambda_{pf}}{E - E_f} V(r') \phi_f(r'). 
\]

The stimulated linewidth of the photoassociated molecule is given by the Fermi–Golden rule expression \( \Gamma_p = 2\pi |A_{pc} h / \hbar | \) and the Feshbach resonance linewidth is \( \Gamma_f = 2\pi |V_{0, fc}^0 h / \hbar | \). Here \( V_{0, fc}^0 = \int dr \phi_f(r) V^*(r) \psi_E^{0, \text{reg}} \) and \( \Lambda_0^0 = \int dr \phi_p(r) A_1(r) \psi_E^{0, \text{reg}} \). The Stark energy shift due to laser coupling of PM state with the continuum is given by \( S_{fc} = \int \text{d}r \text{d}r' \phi_f(r) \Lambda_1^\ast(r) \text{Re}[K(r', r)] A_1(r') \phi_p(r') \). Further, the physics of Feshbach resonance leads us to introduce the parameter

\[
q_f = \frac{\Lambda_{pf} + V_{pf}}{\pi \Lambda_{pc} V_{0, fc}^0} 
\]

which represents an effective continuum-mediated magneto-optical coupling between the two bound states where \( S_{fc} = \int \text{d}r \text{d}r' \phi_f(r) V^*(r) \text{Re}[K(r', r)] V(r') \phi_f(r') \) is the energy shift of the closed-channel bound state due to its coupling with the continuum. Let us introduce a parameter

\[
q_f = \frac{\Lambda_{pf} + V_{pf}}{\pi \Lambda_{pc} V_{0, fc}^0} 
\]

where we call ‘Feshbach asymptmetry parameter’, where \( \Lambda_{pf} = \int dr \phi_p(r) A_1(r) \phi_f(r) \) is the bound–bound Rabi frequency and \( V_{pf} \) is an effective potential between the two bound states due to their interactions with the common ground-state continuum. The expression for \( V_{pf} \) is given in (11). Now writing \( \epsilon = (E - E_f)/|\Gamma_f/2| \) with \( E_f = E_f + S_{fc} \), being the shifted energy of the closed-channel bound state, we can thus express

\[
\tilde{A}_p = \frac{\sqrt{\pi \hbar \Gamma_p^2 / 2}}{h |\Gamma_f/2|} \left( \epsilon + q_f / (\epsilon + 1) \right) \frac{\exp(i\eta_0)}{\Delta_0 + \Gamma_f/2 + D_1} \frac{\exp(i\eta_0)}{\Delta_0 + \Gamma_f/2 + D_1} 
\]

where \( \Delta_0 = [E - (E_p - \hbar \delta)] / |\Gamma_f/2| \) is independent of laser intensity and \( D_1 = (-2S_{pc} + i\hbar \gamma) / \Gamma_f / |\Gamma_f/2| \) is a parameter which is proportional to laser intensity \( I \). In writing the above equation we have assumed that \( \tilde{V}_{fc} \) and \( \tilde{A}_{pc} \) are real quantities. Note that \( q_f \) is independent of laser power since its numerator as well as the denominator is proportional to laser amplitude. Following [23], we can express \( \epsilon \) in terms of applied magnetic field in the form

\[
\epsilon = \frac{E - E_{th} - (E_{th} - E_J)}{\Gamma_f/2} = \frac{E - E_{th} - B - B_0}{\Delta(k a_0)} \Delta(k a_0) 
\]

3
where $E_{th}$ is the threshold of the open channel, $\Delta$ is the Feshbach resonance width, $B_0$ is the resonance magnetic field and $\Delta_{bg}$ is the background scattering length. Here $E - E_{th}$ is the asymptotic collision energy. The energy $E_{th}$ depends on the applied magnetic field due to Zeeman shift of the atomic level. The resonance scattering length is given by $a_{res} = -a_{bg} \Delta/(B - B_0)$.

It is worthwhile to discuss the salient features of our model. The model is capable of accounting for the double resonance comprising of Feshbach and PA resonances. This can be inferred from the expression of PA excitation amplitude $A_p$ given by (13). PA resonance means $\Delta_p = 0$ and Feshbach resonance occurs when $\epsilon = 0$. Putting $\Delta_p = 0$ and $\epsilon = 0$ in (13), we find that $\tilde{A}_p$ remains finite for any value $\gamma$. $\tilde{A}_p$ is expressed in terms of Feshbach asymmetry parameter $q_f$ which governs the quantum interference between two competing optical transition pathways. Here asymmetry refers to the deviation of the spectrum from the standard Lorentzian shape. There are two detuning parameters in the system: one is $\Delta_p$, the frequency offset of the PA laser from the continuum–bound optical transition frequency and the other one is $(E - \tilde{E}_F)$, the magnetic tuning of the continuum threshold with respect to the energy of the bound state (supported by the closed channel) embedded in the continuum. Accordingly, the spectrum can be viewed either as a function of $\Delta_p$ at a fixed $(E - \tilde{E}_F)$ or as a function of $(E - \tilde{E}_F)$ at a fixed $\Delta_p$. When viewed as a function of $\Delta_p$ at a fixed $(E - \tilde{E}_F)$, the spectrum remains Lorentzian with linewidth and line strength largely modified due to the presence of the Feshbach resonance. The modification of linewidth due to MFR is described elsewhere [24]. The asymmetry in the spectral shape arises when the spectrum is plotted as a function of $E - \tilde{E}_F$ or equivalently as a function of the magnetic field. The asymmetry is maximum when $q_f = 0$ and as $q_f \to \infty$, the spectrum becomes the least asymmetric as in the standard Fano profile [1].

Before we discuss our main results, we would like to point out how our mathematical treatment discussed above is related to the recent work of Koznetsova et al [20] who have studied a related model in a different context which is to transfer atoms into ground-state molecules via a two-photon process near a Feshbach resonance. Our approach is to find the real space dressed waveform by solving the time-independent scattering problem by Green’s function method while they have adapted a quantum optics-based approach of finding time-dependent amplitudes of the dressed state by solving coupled differential equations numerically.

3. The results and discussion

We now discuss the characteristic features of our main results. The PA spectrum is given by the PA rate coefficient $K_{PA} = \langle \sigma_{PA} \rangle$, where $\sigma_{PA} = \langle h\gamma | \tilde{A}_p |^2 \rangle/(2\pi k^2)$ is the cross-section for the loss of atoms due to the decay of the excited molecules. Here $\langle \cdot \cdot \cdot \rangle$ implies thermal averaging over the relative velocity $v_{rel} = h k/\mu$. Note that, in the limit $\Gamma_f \to 0$, the PA spectrum reduces to a Lorentzian implying that coupling between the closed-channel bound state and the continuum is essential for the occurrence of Fano interference. When both $\Delta_2$ and $V^{(2)}_{fc}$ go to zero, the spectrum reduces to that of standard PA. For numerical illustrations, we consider a model system of two ground-state ($S_{1/2}$) $^7\text{Li}$ atoms undergoing PA from the ground molecular configuration $^3\Sigma^+_g$ to the vibrational state $v = 83$ of the excited molecular configuration $^3\Sigma^+_u$ which correlates asymptotically to $2S_{1/2} + 2P_{1/2}$-free atoms. The spontaneous linewidth is taken to be $\gamma = 11.7\text{ MHz}$ [25]. The experimental value of shift $S_{pc}$ is reported to be $-1.7 \pm 0.2\text{ MHz/} \text{cm}^2$ [25]. The resonance width is $\Delta = -192.3\text{ Gauss}$ and the background scattering length is $a_{bg} = -24.5 a_0$ ($a_0$ is the Bohr radius) [26]. The Feshbach resonance linewidth $\Gamma_f$ at $10\text{ µK}$ temperature is calculated to be $16.66\text{ MHz}$ using the parameters reported in [27].

Depending on how the PA laser is tuned, we have two cases. In the first case (case-I), the laser is on or near the resonance with free–bound transition but off-resonant with bound–bound transition. In the second case (case-II), it is resonant with bound–bound transition. The PA rate will be maximized at the poles of (13). In case-I, there is only one pole of (13) which depends on laser intensity. The minimum in the spectrum is solely determined by the asymmetry parameter $q_f$ and is independent of laser intensity. We first consider case-I and plot $K_{PA}$ as a function of $B$ for three different values of $\Gamma_p$ in figure 2. For these three different $\Gamma_p$ values we choose three different $q_f$ parameters such that the maximum appears near $B = 736\text{ G}$ [14]. Since the minimum position $B_{min}$ is independent of laser intensity, we also choose three different values of the resonant magnetic field $B_0$ such that $B_{min}$ remains fixed at $710\text{ G}$ for these three $\Gamma_p$ values. In this case there arises the asymmetric Fano profile with one minimum and one maximum. This results from quantum interference between continuum–bound and bound–bound Raman-type transition pathways. This interpretation of the Raman Fano profile is in accordance with the recent experimental observation of two-photon PA by Moal et al [11].

Next we consider case-II in which the PA laser is tuned in resonance with bound–bound rather than continuum–bound transition. In this case we have $h\delta - E_p + \tilde{E}_F = 0$ and so $\Delta_p = \epsilon$. Then $\tilde{A}_p$ will have two maxima given by

$$\epsilon(\epsilon + i) + i\Gamma_f(\epsilon + i) - \Gamma_p(q_f - i)^2 = 0$$

(15)

where $\tilde{\Gamma}_r = \Gamma_r/\Gamma_f$ with $\Gamma_r = \Gamma_r + \gamma$ being the total linewidth, and $\Gamma_p = \Gamma_p/\Gamma_f$. For the sake of comparison, we plot spectra in figure 3 for both cases. Figure 3(a) shows that a single maximum appears in both cases when laser intensities are low. As laser intensity increases, the maximum in case-I disappears while a two-peak structure emerges in case-II as displayed in figure 3(b). We note that the PA rate is lower in case-II in comparison to case-I for the same magnetic field and other parameters except near the two maxima. To further investigate the double-peak structure, we demonstrate spectra for case-II at higher intensities in figure 4 which clearly indicates the nonlinear features of Fano interference. The origin of the two peaks lies in Autler–Townes splitting [2] due to bound–bound resonant coupling when continuum–bound dipole coupling is scanned into a two-photon resonance. At lower intensities, the two peaks can appear on the same side of the Fano minimum. As a result, there can appear another
smaller minimum (we call it Autler–Townes (AT) minimum in order to distinguish it from the Fano minimum) between the two peaks. By comparing figure 3(a) with figure 3(b), we note that the AT minimum at a higher intensity appears near the position where the maximum would have appeared at a lower intensity. The separation between the two maxima increases with increasing laser intensities as shown in figure 4. As one of the peaks crosses the Fano minimum at an increased laser intensity, the AT minimum disappears due to its interference with the much stronger Fano minimum resulting in a two-maximum structure only. The double-maximum structure is particularly prominent in the strong-coupling regime where \( \Gamma_p \) exceeds the spontaneous linewidth \( \gamma \) of the PA molecule.

Recently, Pellegrini and Cote [18] have theoretically obtained double-minimum spectra using the formalism of [19] which is to first diagonalize the part of the Hamiltonian pertaining to the ground-state scattering (continuum interacting with the bound state in closed channel), and then to calculate the optical transition matrix element between this diagonalized state and the excited molecular state by the Fermi–Golden rule. This is similar to linear Fano theory [1] and hence cannot be applied for strong coupling that can further modify the continuum state significantly. In our formalism, we have diagonalized the full Hamiltonian nonperturbatively.

4. Conclusion

The results discussed above clearly demonstrate the linear and nonlinear aspects of Fano interference in weak- and strong-coupling regimes, respectively. The observation of two-minimum and two-maximum structures crucially depends on the precise tuning of PA laser on or near the resonance with bound–bound transition. If the laser field is tuned to get the maximum amount of loss of atoms for a fixed magnetic field, the resulting PA spectrum will mostly correspond to case-I with a single maximum. To explore the nonlinear Fano effect, it is important to know the binding energy of the closed-channel bound state so that the laser can be accurately tuned near the resonance with the bound–bound transition as the magnetic field is varied. Recently, the nonlinear Fano effect was observed in quantum dot [28]. Although Autler–Townes splitting has been recently demonstrated in two-photon PA [10, 11], it is yet to be observed in PA with a single laser beam in the presence of Feshbach resonance. Fano interference may further be explored in photoassociation between heteronuclear atoms such as Na and Cs [29] or K and Rb [16] which have a...
broad magnetic Feshbach resonance and shorter ranged excited potentials.

Acknowledgments

We are thankful to N Bigelow for discussions. This work is supported by the NSF grant no PHYS 0653494.

References

[1] Fano U 1961 Phys. Rev. 124 1866
[2] Rzazewski K and Eberly J H 1981 Phys. Rev. Lett. 47 408
[3] Lambropoulos P and Zoller P 1981 Phys. Rev. A 24 379
   Agarwal G S et al 1982 Phys. Rev. Lett. 48 1164
   Agarwal G S, Haan S L and Cooper J 1984 Phys. Rev. A 29 2552
[4] Harris S E 1989 Phys. Rev. Lett. 62 1033
   Harris S 1997 Phys. Today 50 36
   Lukin M D and Imamoglu A 2001 Nature 413 273
[5] Agarwal G S 1974 Quantum statistical theories of spontaneous emission and their relation to other approaches Springer Tracts in Modern Physics: Quantum Optics (Berlin: Springer)
   Svidzinsky A A, Chang J and Scully M O 2008 Phys. Rev. Lett. 100 160504
[6] Scully M O and Zubairy M S 2003 Science 301 181
[7] Hau L V et al 1999 Nature 397 594
[8] Shi Z et al 2007 Phys. Rev. Lett. 99 240801
[9] Wynar R et al 2000 Science 287 1016
   Winkler K et al 2005 Phys. Rev. Lett. 95 063202
   Ryu C et al 2005 arXiv:cond-mat/0508201
[10] Dumke R et al 2005 Phys. Rev. A 72 041801
[11] Moal S et al 2006 Phys. Rev. Lett. 96 023203
[12] Bohn J L and Julienne P S 1999 Phys. Rev. A 60 414
[13] Bohn J L and Julienne P S 1996 Phys. Rev. A 54 R4637
[14] Junker M et al 2008 Phys. Rev. Lett. 101 000406
[15] Winkler K et al 2007 Phys. Rev. Lett. 98 043201
[16] Ni K K et al 2008 Science 322 231
[17] Mackie M et al 2008 Phys. Rev. Lett. 101 040401
[18] Pellegrini P and Cote R 2009 New J. Phys. 11 055047
[19] Pellegrini P, Gacesa M and Cote R 2008 Phys. Rev. Lett. 101 053201
[20] Kuznetsova E et al 2009 New J. Phys. 11 055028
[21] Tiesinga E, Verhaar B J and Stoof H T C 1993 Phys. Rev. A 47 4114
   Inouye S et al 1998 Nature 392 151
   Courteille Ph et al 1998 Phys. Rev. Lett. 81 69
   Roberts J L et al 1998 Phys. Rev. Lett. 81 5109
[22] Thorsheim H R, Weiner J and Julienne P S 1987 Phys. Rev. Lett. 58 2420
   for reviews on PA, see Weiner J et al 1999 Rev. Mod. Phys. 71 1
   Masnou-Seeuws F and Pillet P 2001 Adv. At. Mol. Phys. 47 53
   Jones K M et al 2006 Rev. Mod. Phys. 78 483
[23] Moerdijk A J, Verhaar B J and Axellson A 1995 Phys. Rev. A 51 4852
[24] Deb B and Rakshit A 2009 J. Phys. B: At. Mol. Opt. Phys. 42 195202
[25] Prodan I D et al 2003 Phys. Rev. Lett. 91 080402
[26] Pollack S E et al 2009 Phys. Rev. Lett. 102 090402
[27] Chin C et al 2008 LANL arXiv 0812.1496
[28] Kroner M et al 2008 Nature 451 311
[29] Shaffer J P, Chalupczak W and Bigelow N P 1999 Phys. Rev. Lett. 82 1124
   Haimberger C et al 2004 Phys. Rev. A 70 021402