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Optical control of magnetic Feshbach resonances in Bose gases

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ABSTRACT

We theoretically investigate optical control of magnetic Feshbach resonance in Bose gases with two optical fields. The two optical fields couple two ground states through an excited state. Compared with the usual single-optical scheme, two optical fields can greatly suppress the inelastic loss resulting from spontaneous emission by the destructive quantum interference. Using the mean field theory, the analytical formula of the scattering length is obtained. The results show that the scattering length can be modified in a large range by changing the Rabi frequency or the optical field frequency. The strong atom–molecule interaction has obvious effect on the scattering length.

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1. Introduction

Ultracold atoms and molecules have attracted much attention of researchers due to their application in precision measurement [1], the quantum computation [2], and chemical reaction [3]. Control of atomic interactions has made an explosive progress in quantum gases for investigating few- and many-body quantum systems.

The magnetic Feshbach resonance (MFR) provides a useful tool to control the interaction between atoms in ultracold quantum gases [4–8]. However, magnetic modulation is limited in current experiments of MFR to relatively longer time scales and lower spatial resolution. Optical Feshbach resonance (OFR) is expected to be a new and powerful control with higher spatial and temporal resolutions [9]. OFR has been studied in thermal degenerate gases of Na [10] and ^{87}Rb [11], but for alkali atoms it may suffer from rapid loss of atoms due to light-induced inelastic collisions. The loss is characterized by a two-body rate coefficient K_2 with an estimated value of 10^{-10} cm^3 s^{-1} [12]. Typical density of the order of 10^{14} cm^{-3} in experiment results in lifetime of the order of 100 μs [12], which is too short for many applications. OFRs in ultracold alkaline–earth atomic collisions may overcome the loss problem [13]. OFR for control of higher partial wave, such as p-wave scattering of ^{171}Yb, has been suggested [14] and demonstrated [15]. The Bose gas ^{88}Sr with a narrow intercombination line has been studied by OFR, using a laser with frequency tuned away from resonance [16].

Recently, the combination of MFR and OFR has aroused a great interest both in theory and experiment. Bauer et al. demonstrated that a single-color laser can not only noticeably shift an MFR in ^{87}Rb [12,17], but also induce considerably low rates of particle loss than an OFR. Wu et al. proposed a way for optical control of MFR in Fermi gases using two optical fields [18,19]. This dark-state optical method can be used to control the interaction strength near an MFR, and suppress spontaneous emission by quantum interference. Deb et al. suggested that quantum interference between magnetic and optical Feshbach resonances could suppress inelastic scattering and enhance elastic scattering cross section [20]. Furthermore, this method can be used to manipulate not only the spherically symmetric s-wave interaction, but also the anisotropic higher partial-wave interaction.

A many-body approach was applied to describe MFRs by Timmermans et al. [21]. In the present work, this many-body approach is extended to describe the optical control of MFRs in Bose gas. Using the mean field theory, the analytical formula of the scattering length is obtained. The atom–molecule interaction, which has been not considered in previous studies [12,17–20], is taken into account and its effects on the scattering length are discussed.

2. Theoretical approach

Fig. 1 illustrates a scheme for a pair of atoms in the hyperfine state, which undergoes an s-wave collision in the ground electronic state (open channel). The hyperfine interaction can couple the scattering continuum state $|a\rangle \otimes |a\rangle$ of the open channel to a bound vibrational state $|g_1\rangle$ in the closed channel. The energy of colliding atomic pair can be modulated close to state $|g_1\rangle$ by a magnetic field, resulting in an MFR. Another ground state which is not coupled to the open channel is a certain vibrational state. The optical

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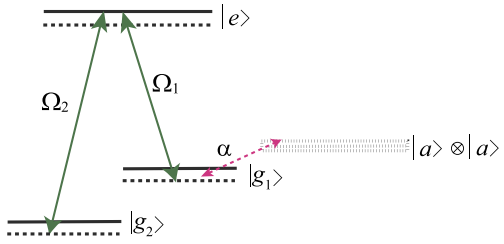

Fig. 1. A scheme for optical control of a Feshbach resonance with two optical fields. The two optical fields with frequencies ω_1 and ω_1 and Rabi frequencies Ω_1 and Ω_2 couple the ground states $|g_1\rangle$ and $|g_2\rangle$ to the excited state $|e\rangle$, respectively. α denotes the coupling between the incoming atomic pair state $|a\rangle\otimes|a\rangle$ and $|g_1\rangle$, which is responsible for a magnetically induced Feshbach resonance.

fields Ω_1 and Ω_2 couple the ground states $|g_1\rangle$ and $|g_2\rangle$ to the excited electronic state $|e\rangle$, respectively. The Hamiltonian of colliding system is written as [17]

$$\begin{aligned}\hat{H} = &\int d^3\mathbf{x}\sum_j \hat{\Psi}_j^\dagger(\mathbf{x})H_j(\mathbf{x})\hat{\Psi}_j(\mathbf{x})\\ &+\frac{1}{2}\int d^3\mathbf{x}\sum_{i,j}\hat{\Psi}_i^\dagger(\mathbf{x})\hat{\Psi}_j^\dagger(\mathbf{x})U_{ij}\hat{\Psi}_j(\mathbf{x})\hat{\Psi}_i(\mathbf{x})\\ &+\hbar\int d^3\mathbf{x}_1 d^3\mathbf{x}_2\big[\hat{\Psi}_1^\dagger(\mathbf{X}_{12})\alpha(\mathbf{x}_{12})\hat{\Psi}_p(\mathbf{x}_1,\mathbf{x}_2)+\text{H.c.}\big]+\hat{H}_{\text{int}}\end{aligned} \tag{1}$$

where the optical field–system interaction Hamiltonian is given by

$$\begin{aligned}\hat{H}_{\text{int}} = -\hbar\int d^3\mathbf{x}\big[&\cos(\omega_1 t)\hat{\Psi}_e^\dagger(\mathbf{x})\Omega_1\hat{\Psi}_1(\mathbf{x})\\ &+\cos(\omega_2 t)\hat{\Psi}_e^\dagger(\mathbf{x})\Omega_2\hat{\Psi}_2(\mathbf{x})+\text{H.c.}\big].\end{aligned} \tag{2}$$

Here, H_j $(j=a,1,2,e)$ is the single-particle (atom or molecule) Hamiltonian, $\hat{\Psi}_j(\mathbf{x})$ $(\hat{\Psi}_j^\dagger(\mathbf{x}))$ is the bosonic annihilation (creation) field operator. The relative coordinate $\mathbf{x_{12}}=\mathbf{x_1}-\mathbf{x_2}$ and the center-of-mass coordinate $\mathbf{X_{12}}=(\mathbf{x_1}+\mathbf{x_2})/2$. $\hat{\Psi}_p(\mathbf{x_1},\mathbf{x_2})=\hat{\Psi}_a(\mathbf{x_1})\hat{\Psi}_a(\mathbf{x_2})$ denotes the annihilation field operator of an atomic pair. $U_i=U_{ii}=4\pi\hbar^2 a_i/m_i$ represents the two-body interaction strength between the same species, and $U_{ij}=U_{ji}=2\pi\hbar^2 a_{ij}/m_{ij}$ (for $i\neq j$) denote the interaction strengths between different species, where a_i and a_{ij} are the corresponding s-wave scattering lengths far away from Feshbach resonance. $m_j=m(2-\delta_{aj})$ is the mass of a particle in the jth state with δ_{aj} being the Kronecker symbol and m the atomic mass, and $m_{ij}=m_i m_j/(m_i+m_j)$ the reduced mass. In this paper, the molecule–molecule interactions are neglected. Only the atom–atom interactions and the interactions between atom and molecule in the ground molecular states are considered. $\alpha(\mathbf{x}_{12})$ denotes the atom–molecule coupling relevant to the MFR. $\Omega_{1(2)}=\langle e|d|g_{1(2)}\rangle\varepsilon_{1(2)}/\hbar$ is the Rabi frequency corresponding to the transition $|g_1\rangle\to|e\rangle$ $(|g_2\rangle\to|e\rangle)$.

If the incoming atoms are slow enough, $\alpha(\mathbf{x}_{12})$ approximates to a constant α [17]. $\hat{H}$ in Eq. (1) is computed by $\hat{H}=\int d^3\mathbf{x}\,\mathcal{H}(\mathbf{x})$ with the Hamiltonian density [21]

$$\begin{aligned}\mathcal{H}(\mathbf{x}) = &\sum_j\hat{\Psi}_j^\dagger(\mathbf{x})H_j(\mathbf{x})\hat{\Psi}_j(\mathbf{x})+\frac{1}{2}U_a\hat{\Psi}_p^\dagger(\mathbf{x})\hat{\Psi}_p(\mathbf{x})\\ &+\frac{1}{2}\sum_{j=1,2}U_{aj}\big(\hat{\Psi}_a^\dagger(\mathbf{x})\hat{\Psi}_j^\dagger(\mathbf{x})\hat{\Psi}_j(\mathbf{x})\hat{\Psi}_a(\mathbf{x})\\ &+\hat{\Psi}_j^\dagger(\mathbf{x})\hat{\Psi}_a^\dagger(\mathbf{x})\hat{\Psi}_a(\mathbf{x})\hat{\Psi}_j(\mathbf{x})\big)+\hbar\big[\hat{\Psi}_1^\dagger\alpha\hat{\Psi}_p+\text{H.c.}\big]\\ &+\mathcal{H}_{\text{int}}(\mathbf{x}),\end{aligned} \tag{3}$$

where

$$\begin{aligned}\mathcal{H}_{\text{int}}(\mathbf{x}) = -\hbar\big[&\cos(\omega_1 t)\hat{\Psi}_e^\dagger(\mathbf{x})\Omega_1\hat{\Psi}_1(\mathbf{x})+\cos(\omega_2 t)\hat{\Psi}_e^\dagger(\mathbf{x})\Omega_2\hat{\Psi}_2(\mathbf{x})\\ &+\text{H.c.}\big].\end{aligned} \tag{4}$$

The single-particle Hamiltonian containing a kinetic operator and internal energy E_j^{int} is given by

$$H_j(\mathbf{x}) = -\frac{\hbar^2}{2m_j}\nabla^2+E_j^{\text{int}} \tag{5}$$

with $j=(a,1,2,e)$. The internal energy E_j^{int} depends nonlinearly on the magnetic field B. Near the pole of the unshifted Feshbach resonance B_0, this dependence is linear and expressed as

$$E_j^{\text{int}} = -\mu_j(B-B_0)+\hbar\omega_{e1}\delta_{ej}+\hbar\omega_{21}\delta_{2j}, \tag{6}$$

where μ_j is magnetic dipole moment. At $B=B_0$, the $|a\rangle$ and $|g_1\rangle$ states are degenerate, and the energies of the $|g_2\rangle$ and $|e\rangle$ states are $\hbar\omega_{21}$ and $\hbar\omega_{e1}$, respectively.

We assume that the population in each state is Bose condensed and can be described by a mean field $\psi_j(\mathbf{x})=\langle\hat{\Psi}_j(\mathbf{x})\rangle$ for a homogeneous system. By taking the expectation value of the Heisenberg equation of motion $i\hbar\,d\hat{\Psi}_j/dt=[\hat{\Psi}_j,\hat{H}]$ and approximating the field operators as uncorrelated, one can obtain [21]

$$i\frac{d\psi_a}{dt} = \frac{E_a^{\text{int}}}{\hbar}\psi_a+\frac{1}{\hbar}\sum_{j=a,1,2}U_{aj}|\psi_j|^2\psi_a+2\alpha^*\psi_a^*\psi_1, \tag{7a}$$

$$i\frac{d\psi_1}{dt} = \frac{E_1^{\text{int}}}{\hbar}\psi_1+\frac{1}{\hbar}U_{a1}|\psi_a|^2\psi_1+\alpha\psi_a^2-\Omega_1^*\psi_e\cos(\omega_1 t), \tag{7b}$$

$$i\frac{d\psi_2}{dt} = \frac{E_2^{\text{int}}}{\hbar}\psi_2+\frac{1}{\hbar}U_{a2}|\psi_a|^2\psi_2-\Omega_2^*\psi_e\cos(\omega_2 t), \tag{7c}$$

$$i\frac{d\psi_e}{dt} = \frac{E_e^{\text{int}}}{\hbar}\psi_e-\Omega_1^*\psi_1\cos(\omega_1 t)-\Omega_2^*\psi_2\cos(\omega_2 t). \tag{7d}$$

By taking $\psi_j=b_j e^{-i[(2-\delta_{aj})E_a^{\text{int}}/\hbar+\omega_1\delta_{ej}-(\omega_2-\omega_1)\delta_{2j}]t}$ and using the rotating wave approximation, Eq. (7) can be reduced to

$$i\frac{db_a}{dt} = \frac{1}{\hbar}\sum_{j=a,1,2}U_{aj}|b_j|^2 b_a+2\alpha^* b_a^* b_1, \tag{8a}$$

$$i\frac{db_1}{dt} = \Delta_1 b_1+\frac{1}{\hbar}U_{a1}|b_a|^2 b_1+\alpha b_a^2-\frac{1}{2}\Omega_1^* b_e, \tag{8b}$$

$$i\frac{db_2}{dt} = \Delta_2 b_2+\frac{1}{\hbar}U_{a2}|b_a|^2 b_2-\frac{1}{2}\Omega_2^* b_e, \tag{8c}$$

$$i\frac{db_e}{dt} = -\left(\Delta_e+\frac{i\gamma_e}{2}\right)b_e-\frac{1}{2}\Omega_1 b_1-\frac{1}{2}\Omega_2 b_2, \tag{8d}$$

where $\Delta_1=\mu_{a1}(B-B_0)/\hbar$, the single-photon detuning $\Delta_e=\omega_1-\mu_{ae}(B-B_0)/\hbar$ and the two-photon detuning $\Delta_2=(\omega_2-\omega_1)-\mu_{2a}(B-B_0)/\hbar$ with $\mu_{a1}=\mu_a-\mu_1$ and $\mu_{2a}=\mu_2-\mu_a$. The radiative decay rate of the molecule in the excited state is $\gamma_e=2/\tau_{\text{spont}}$, where τ_{spont} is the atom spontaneous lifetime. γ_e is used to describe the photoassociation loss rates near an MFR.

It is assumed that all population is initially in state $|a\rangle$ and the populations in states $|g_1\rangle$, $|g_2\rangle$ and $|e\rangle$ are approximately zero [22]. This is a good approximation, if the angular frequency $\alpha b_a\ll\Omega_1$ and $db_e/dt\ll\gamma_e b_e/2$. This condition is always satisfied in the low-density limit, but for a very broad Feshbach resonance it may be difficult to satisfy this condition experimentally.

The adiabatic elimination is done by setting $db_1/dt=db_2/dt=db_e/dt=0$, yielding

$$i\frac{db_a}{dt} = \sum_{j=1}^{6}Q_{2j}|b_a|^{2j}b_a \tag{9}$$

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