Optical Control of Exchange Interaction and Kondo Temperature in cold Atom Gas

Igor Kuzmenko¹, Tanya Kuzmenko¹ and Yshai Avishai¹,²,³

¹Department of Physics, Ben-Gurion University of the Negev, Beer-Sheva, Israel
²New York University and the NYU-ECNU Institute of Physics at NYU Shanghai, 3663 Zhongshan Road North, Shanghai, 200062, China
³Yukawa Institute for Theoretical Physics, Kyoto, Japan

(Dated: September 24, 2018)

PACS numbers: 37.10.Jk, 31.15.vn, 33.15.Kr

The relevance of magnetic impurity problems in cold atom systems depends crucially on the nature of the exchange interaction between itinerant fermionic atoms and localized impurity atoms. In particular, Kondo physics occurs only if the exchange interaction is anti-ferromagnetic, and strong enough to yield high enough Kondo temperature \( (T_K/T_F \geq 0.1) \). Focusing, as an example, on the experimentally accessible system of ultra-cold \(^{173}\text{Yb}\) atoms, it is shown that the sign and strength of an exchange interaction between an itinerant \(^{173}\text{Yb}(^3S_0)\) atom and a trapped \(^{173}\text{Yb}(^3P_0)\) atom can be optically controlled. Explicitly, as the light intensity increases (from zero), the exchange interaction changes from ferromagnetic to anti-ferromagnetic. When the light intensity is such that the system approaches a singlet Feshbach resonance (from below), the singlet scattering length \( a_S \) is large and negative, and the Kondo temperature increases sharply.

Introduction: Controlling interaction between cold atoms is a godsend, as it turns atomic systems capable of demonstrating new phenomena that cannot be otherwise accessed within solid state physics proper. In a series of theoretical and experimental investigations, it has been established that, as far as potential scattering is concerned, the strength of atomic interactions can be tuned by laser beams. The prime object of these studies is to achieve an optical Feshbach resonance and thereby to obtain a Bose-Einstein condensate in cold bosonic atom systems.

On the other hand, the feasibility of controlling the strength and the sign of exchange interaction between atoms is much less studied. Its importance has been recognized recently, in the quest for studying the Kondo effect in cold atom systems. The physics exposed in the study of magnetic impurities when the itinerant (fermionic) atoms have spin \( F \geq \frac{3}{2} \) is very rich, touching upon exotic phenomena such as over-screening, realization of the Coqblin-Schrieffer model, multipolar Kondo effect and others. Recently, it has been shown that exchange interaction can be controlled by the technique of confined-induced-resonance (CIR).

The goal of the present study is to show that exchange interaction between fermionic atoms can be optically controlled. As an experimentally feasible example we consider \(^{173}\text{Yb}\) atoms. Those in the ground state \(^1S_0\) are itinerant, (forming a degenerate Fermi gas confined in a shallow square well potential), whereas those in the long lived excited state \(^3P_0\) are trapped in a state-dependent optical potential and serve as dilute concentration of localized magnetic impurities. Both the ground \(^1S_0\) and excited \(^3P_0\) state atoms have spin \( F = \frac{5}{2} \) (that is the nuclear spin). In this case, the Coqblin-Schrieffer model with SU(6) symmetry is realizable due to a unique exchange mechanism. It is shown that by applying laser beams on the atomic gas, the exchange interaction between \(^{173}\text{Yb}(^1S_0)\) and \(^{173}\text{Yb}(^3P_0)\) atoms can be controlled, both in sign and magnitude. In particular, with increasing the light intensity, the exchange interaction changes from ferromagnetic to anti-ferromagnetic, that is a prerequisite for occurrence of the Kondo physics.

Description of the system: Consider a 3D gas of \(^{173}\text{Yb}\) atoms confined in a shallow square well potential. Most of the atoms remain in the ground state \(^1S_0\) and form a Fermi sea due to its half integer nuclear spin \( I = \frac{5}{2} \) (see green area in Fig. 1). However, following a coherent excitation via the clock transition, a few atoms occupy the long-lived excited state (electronic configuration \(^3P_0\) ). These excited atoms can be trapped in a state-dependent optical lattice potential as schematically displayed in Fig. 1 (red circles), and can be regarded as dilute concentration of localized impurities.

![Fig. 1: (Color online) Illustration of a magnetic impurity setup for \(^{173}\text{Yb}\) atoms (schematic). Atoms in the ground-state \(^1S_0\) form a Fermi sea (Fermi energy \( \epsilon_F \)), while atoms in the excited-state \(^3P_0\) are trapped in an optical potential and form a dilute concentration of localized magnetic impurities. \( \epsilon_{e,g} \) is the excitation energy of the \(^3P_0\) state.](image-url)
The energy dispersion and density of states (DOS) for the (weakly confined) Yb\(^{(1S_0)}\) atoms read,

\[ \epsilon_k = \frac{\hbar^2 k^2}{2M}, \quad \rho(\epsilon) = \frac{M k_e}{2\pi^3 \hbar^2}. \]  

(1)

Here \( M \) is the atomic mass, \( k = |\mathbf{k}| \), \( \mathbf{k} = (k_x, k_y, k_z) \), with \( k_{\alpha=x,y,z} = \frac{\pi n_{\alpha}}{a_0} \), in which \( \{n_\alpha\} \) are integers and \( k_e = \sqrt{2M\epsilon/h} \).

As for the impurities, an Yb\(^{(3P_0)}\) atom at position \( R = (X_1, X_2, X_3) \) is trapped by an optical potential,

\[ V_e(R) = -V_e^{(0)} \sum \sin^2 (k_e X_\alpha), \]  

(2)

where \( k_e \) is the wave number of the laser light. The lowest energy level \( \epsilon_{\text{imp}} \) of the Yb\(^{(3P_0)}\) atom is deep and close to the minimum of \( V_e(R) \), hence it can be approximated harmonic potential at lowest energy,

\[ \epsilon_{\text{imp}} = \frac{3}{2} \hbar \Omega_e, \]  

(3)

wherein the harmonic frequency and length are,

\[ \hbar \Omega_e = 2\sqrt{E_e V_e^{(0)}}, \quad k_e a_e = \left( \frac{E_e}{V_e^{(0)}} \right)^{1/4}, \]  

(4)

with recoil energy \( E_e = \frac{\hbar^2 k_e^2}{2M} \).

**Exchange interaction** between the Yb\(^{(1S_0)}\) and Yb\(^{(3P_0)}\) atoms is described in details in Ref. [19]. Here we briefly discuss it and explain how it can be optically controlled. Consider an Yb atom as a doubly ionized closed shell rigid ion and two valence electrons. The ground state \( ^1S_0 \) electron configuration is \( 6s^2 \), whereas that of the excited state \( ^3P_0 \) is \( 6s6p \). The excitation energy is

\[ \epsilon_{e,g} = \epsilon_e - \epsilon_g = 2.14349 \text{ eV}. \]  

(5)

The positions of the ion core and the outer electrons are respectively specified by vectors \( \mathbf{R} \), \( \mathbf{r}_s \), and \( \mathbf{r}_p \). In the adiabatic (Born-Oppenheimer) approximation (which is well substantiated in atomic physics), the wave function of a single Yb atom is expressed as a product of the wave functions of the rigid ion core and of the valence electrons. The former is considered as a point particle of mass \( M \). When one valence electron virtually tunnels from Yb\((6s^2)\) to Yb\((6s6p)\) we get an ionized Yb\(^{+}(6s)\) atom and a charged Yb\(^{-(6s^26p)}\) atom. The excitation energy of this state is

\[ \Delta \epsilon = \epsilon_{\text{ion}} - \epsilon_{\text{ea}} - \epsilon_{e,g} = 4.4107 \text{ eV}, \]  

(6)

gives rise to anti-ferromagnetic exchange interaction between them [19].

\[ V_{\text{exch}}^{(\text{bare})}(R) = -\frac{2}{\Delta \epsilon} t_{s,\text{bare}}(R) t_{p,\text{bare}}(R), \]  

(7)

where \( R = |\mathbf{R}_g - \mathbf{R}_e| \) is the distance between the atoms and \( \Delta \epsilon \) is given by eq. [6] [negative sign of \( V_{\text{exch}}^{(\text{bare})}(R) \) implies an anti-ferromagnetic exchange]. \( t_{s,\text{bare}}(R) \) and \( t_{p,\text{bare}}(R) \) are the tunnelling rates for the 6s and 6p electrons [19], given by,

\[ t_{\mu,\text{bare}}(R) = -\int d^3r \tilde{U}(r) \psi_\mu^*(r) \psi_\mu(|r - \mathbf{R}|), \]  

(8)

where \( \mu = s, p \), \( \mathbf{R} = \mathbf{R}_g - \mathbf{R}_e \), and \( \psi_\mu(r) \) is a radial single-electron wave function of the 6s or 6p electron. The potential

\[ \tilde{U}(r) = U_t(r) + e^2 \int_0^r \left| \psi_s(r') \right|^2 \left( 1 + \frac{r'}{r} \right) r' dr', \]  

(9)

specifies the interaction between the 6s or 6p electron with the positively charged ion Yb\(^+(6s)\), and \( U_t(r) \) is the potential of interaction of the outer electron with the double ionized rigid ion Yb\(^{2+}\). For the wave function \( \psi_\mu(r) \), we use the expression [19],

\[ \psi_\mu(r) = \frac{2A \kappa_\mu^3}{\pi \Gamma \left( \frac{3+2}{3+2} \right)} 2^{2+3} e^{-\kappa_\mu r}, \]  

(10)

where \( A = 0.84095 \),

\[ \beta_s = 0.67680, \quad \kappa_s = 1.2812 \text{ Å}^{-1}, \]

\[ \beta_p = 0.54967, \quad \kappa_p = 1.0387 \text{ Å}^{-1}. \]

When \( r \) exceeds the radius of the Yb\(^{2+}\) ion (which is much shorter than the atomic radius), we can write \( U_t(r) = -\frac{2e^2}{r} \).

**Van der Waals interaction:** The van der Waals interaction between the Yb atoms is modelled here by the Lennard-Jones potential [27],

\[ W^{(\text{bare})}(R) = \frac{C_6}{R^6} \left\{ \frac{\sigma^6}{R^6} - 1 \right\} - \frac{C_8}{R^8}, \]  

(11)

with \( C_6 = 1.9317 \cdot 10^3 E_h a_0^6, \) \( C_8 = 1.94961 \cdot 10^5 E_h a_0^8 \) and \( \sigma = 9.0109362 a_B \), where \( E_h = 27.211 \text{ eV} \) is the Hartree energy and \( a_B = 0.52918 \text{ Å} \) is the Bohr radius.

**Light-assisted interaction:** Controlling the strength and sign of the exchange interaction is achieved by subjecting the mixture of Yb atoms to a laser beam of frequency \( \omega_0 \) that is tuned to be close to the resonant frequency \( \omega_{\text{res}} = \Delta \epsilon / \hbar \). Recall that \( \Delta \epsilon \) is the energy difference between the two neutral atoms and the two ions. Concretely, \( \omega_{\text{res}} = 35574.7 \text{ cm}^{-1} \). The laser light
induces an additional potential $W^{(\text{ind})}(R)$ and exchange interactions $V_{\text{exch}}^{(\text{ind})}(R)$ between the neutral atoms,

$$W^{(\text{ind})}(R) = -\frac{1}{\hbar \omega_{\text{Rab}}} \left\{ r^2_{\text{s,ind}}(R) + r^2_{\text{p,ind}}(R) \right\}, \quad (12)$$

$$V_{\text{exch}}^{(\text{ind})}(R) = -\frac{2}{\hbar \omega_{\text{Rab}}} t_{s,\text{ind}}(R) t_{p,\text{ind}}(R), \quad (13)$$

where the Rabi (detuning) frequency is

$$\omega_{\text{Rab}} = \omega_{\text{res}} - \omega_0.$$  

The tunnelling rates $t_{\mu,\text{ind}}(R)$ ($\mu = s, p$) are,

$$t_{\mu,\text{ind}}(R) = \frac{e v_0 E_0}{4 \pi \omega_0} \mathcal{F}_\mu(R),$$

where $v_0 = \frac{\hbar}{m_e \omega_0} = 0.00729735c$, $E_0$ is the amplitude of the laser’s electric field and the dimensionless functions $\mathcal{F}_\mu(R)$ are

$$\mathcal{F}_\mu(R) = a B \int d^3 r \, \psi^*_{\mu}(|r - R|) \frac{\partial \psi_{\mu}(r)}{\partial z}. \quad (14)$$

The axis $z$ is chosen parallel to the vector $R = R_s - R_c$. It is useful to rewrite the (laser induced) potential and exchange contributions [$W^{(\text{ind})}(R)$ and $V_{\text{exch}}^{(\text{ind})}(R)$] in the compact form,

$$W^{(\text{ind})}(R) = -\frac{V_0}{2} \left\{ F^2_s(R) + F^2_p(R) \right\}, \quad (15)$$

$$V_{\text{exch}}^{(\text{ind})}(R) = -V_0 \, F_s(R) \, F_p(R), \quad (16)$$

where the coupling $V_0$ is given by,

$$V_0 = \frac{2}{\hbar \omega_{\text{Rab}}} \left( \frac{e E_0 v_0}{4 \pi \omega_0} \right)^2. \quad (17)$$

**Scattering lengths:** The van der Waals and exchange interactions yield “singlet” and “triplet” scattering lengths [19, 27, 28],

$$a_\nu = \tilde{a} \left\{ 1 - \tan \left( \Phi_\nu - \frac{\pi}{8} \right) \right\}, \quad (18)$$

where $\nu = S, T$ for the quantum states with antisymmetric (“singlet”) and symmetric (“triplet”) two-particle spin wave functions, and

$$\tilde{a} = \frac{1}{2^{7/2}} \frac{\Gamma \left( \frac{7}{4} \right)}{\Gamma \left( \frac{5}{4} \right)} \left( \frac{MC_6}{\hbar^2} \right)^{1/4} = 39.73 \, \text{Å}. \quad (19)$$

The parameters $\Phi_\nu$ are,

$$\Phi_\nu = \int_{R_\nu} K_\nu(R) \, dR, \quad (20)$$

where

$$K_\nu(R) = \frac{1}{\hbar} \sqrt{-M \left[ W(R) + \eta_\nu V_{\text{exch}}(R) \right]} \quad (21)$$

In the above equation $W(R) = W^{(\text{bare})}(R) + W^{(\text{ind})}(R)$, $V_{\text{exch}}(R) = V_{\text{exch}}^{(\text{bare})}(R) + V_{\text{exch}}^{(\text{ind})}(R)$, $\eta_S = 1$ and $\eta_T = -1$. $R_\nu$ is the solution of equation $W(R_\nu) + \eta_\nu V_{\text{exch}}(R_\nu) = 0$.

When the intensity of the laser beam with frequency $\omega_0$ vanishes [i.e., when $V_0 \to 0$], the scattering lengths are [19]

$$a_S = 1149.9 \, \text{Å}, \quad a_T = 115.9 \, \text{Å}, \quad (22)$$

which agree well with the experimental results [29].

Evaluating the scattering lengths for nonzero $V_0$, requires calculation of $V_{\text{exch}}^{(\text{ind})}(R)$ and $W^{(\text{ind})}(R)$, eqs. (16) and (17). Substituting them into eq. (20) yields the quantity $\Phi_{\nu=S,T}$ as functions of $V_0$. This is carried out numerically, after which eq. (18) is employed to find the scattering lengths.

**FIG. 2:** (Color online) Scattering lengths $a_S$ [blue curve] and $a_T$ [red curve] as functions of $V_0$.

The scattering lengths [18] calculated numerically are displayed in Fig. 2 as functions of $V_0$. It is seen that in the interval $V_0 < 26$ meV, there are two values of $V_0$ where $a_S = a_T$: $V_{c,1} = 1.93$ meV and $V_{c,3} = 15.53$ meV. In addition, there are two values of $V_0$ where $a_S$ is singular: $V_{c,2} = 11.15$ meV and $V_{c,4} = 22.43$ meV. On the other hand, $a_T$ changes slowly within the interval $V_0 < 26$ meV.

**Kondo Hamiltonian:** Once the “singlet” and “triplet” scattering lengths are known, it is possible to construct the effective Kondo Hamiltonian, with explicit expressions for potential and exchange coupling constants denoted below as $g_{\text{pot}}$ and $g_{\text{exch}}$ respectively. Let us consider a degenerate Fermi gas of Yb($^1S_0$) atoms, and one Yb($^3P_0$) atom localized at $R_c = 0$ which plays a role of an impurity. Then the Kondo Hamiltonian is,

$$H_K = g_{\text{pot}} \sum_{k,k'} \sum_m c_{k',m}^\dagger c_{k,m} +$$

$$+ g_{\text{exch}} \sum_{k,k',m \neq m'} \left\{ X^{m,m'} c_{k',m'}^\dagger c_{k,m} +$$

$$+ g_{\text{exch}} \sum_{k,k'} \sum_m Z^{m,m'} c_{k',m'}^\dagger c_{k,m}. \quad (23)$$

Here $c_{k,m}$ and $c_{k,m}^\dagger$ are annihilation and creation operators of itinerant atom with wave vector $k$ and mag-
netic quantum number \( m \). \( X^{m,m'} = |m\rangle\langle m'| \) are Hubbard operators of the localized impurity, the ket \( |m\rangle \) describes the impurity with magnetic quantum number \( m \). \( Z^{m,m} = X^{m,m} + \frac{1}{6} \). \( g_{\text{pot}} \) and \( g_{\text{exch}} \) are effective couplings of the potential and exchange interactions. Here \( g_{\text{exch}} > 0 \) denotes antiferromagnetic exchange interaction. The laser induced scattering lengths due to the short range interaction \([23]\) are,

\[
a_{\nu=S,T} = \frac{M}{4\pi\hbar^2} \left( g_{\text{pot}} + \alpha_{\nu} g_{\text{exch}} \right), \tag{24}
\]

where \( \alpha_S = -\frac{7}{4} \) and \( \alpha_T = \frac{5}{4} \). Then the couplings \( g_{\text{pot}} \) and \( g_{\text{exch}} \) of the potential and exchange interactions are expressed in terms of the scattering lengths as,

\[
g_{\text{pot}} = \frac{\pi\hbar^2}{3M} \left\{ 5a_S + 7a_T \right\}, \tag{25}
\]

\[
g_{\text{exch}} = \frac{2\pi\hbar^2}{M} \left\{ a_T - a_S \right\}. \tag{26}
\]

Eq. \([20]\) shows that when \( a_T > a_S \), the exchange interaction is anti-ferromagnetic and the Hamiltonian \([23]\) gives rise to Kondo effect. When \( a_T < a_S \), the exchange interaction is ferromagnetic and there is no Kondo effect.

**Kondo temperature**: When the exchange interaction \([23]\) is anti-ferromagnetic (corresponding to regions in Fig. 2 where \( a_T > a_S \)), the Kondo temperature is \([10]\),

\[
T_K = D_0 \exp \left( -\frac{1}{6 g_{\text{exch}} \rho(\epsilon_F)} \right). \tag{27}
\]

Hereafter we assume that \( 2\hbar\Omega_e < \epsilon_F \) [where \( \epsilon_F \) is the Fermi energy and the harmonic frequency \( \Omega_e \) is given by eq. \([4]\), and therefore \( D_0 = 2\hbar\Omega_e \) plays a role of ultraviolet cutoff of the Kondo theory. The constant \( \rho(\epsilon_F) \) is the DOS \([4]\) at the Fermi energy. A simple expression relating \( T_K \) to the scattering lengths (valid for \( a_T > a_S \)) then reads,

\[
T_K = D_0 \exp \left( -\frac{\pi}{6k_F(a_T - a_S)} \right), \tag{28}
\]

where \( a_T \) and \( a_S \) as functions of \( V_0 \) are shown in Fig. 2.

The Kondo temperature \([25]\) calculated numerically is shown in Fig. 3 (blue curves). It is seen that the Kondo temperature sharply increases with decreasing \( a_S \) [see Fig. 2]. The intervals of \( V_0 \) where \( a_S > a_T \) and there is no Kondo effect are marked by red.

**Conclusions**: Tuning interaction strength in quantum impurity problems in cold atom systems cannot rely on the application of an external magnetic field (for driving Feshbach resonance), because it is detrimental for the Kondo effect. Hence, employing optical toolbox for controlling interaction strength in cold atom systems is a proper substitute. But so far it has been applied in numerous works mainly for studying bosonic systems. The quest for studying quantum (magnetic) impurity problems and Kondo physics requires a novel kind of controlling the exchange interaction. More concretely, it involves a subtle tuning of “singlet” and “triplet” scattering lengths, and identifying the conditions wherein \( a_S < a_T \), for which the exchange interaction is antiferromagnetic.

This objective has been achieved here for an experimentally representative system. The feasibility to construct optically tuneable exchange interaction between itinerant \( ^{173}\text{Yb}(^1S_0) \) atoms and a localized \( ^{173}\text{Yb}(^{3}P_0) \) impurity has been analyzed. “Singlet” and “triplet” scattering lengths as a function of \( V_0 \) (which is proportional to the light intensity) are explicitly calculated and the regions where \( a_S < a_T \) in which the exchange interaction is antiferromagnetic are identified. With increasing intensity of light (from zero), the exchange interaction changes both in magnitude and in sign. The Kondo Hamiltonian is then constructed and the Kondo temperature is calculated in the intervals of \( V_0 \) where the exchange interaction is anti-ferromagnetic. It is shown that \( T_K \) increases sharply before reaching an optical Feshbach resonance where the singlet scattering length approaches \(-\infty\) \([21]\).

Acknowledgement We thank the authors of Refs.\([21–24]\) for drawing our attention to the CIR method. Prof. Y. Band is to be acknowledged for pointing out the relevance of spontaneous emission. This research is supported in part by an Israel Science Foundation grant 400/12.

---

[1] S. E. Pollack, D. Dries, M. Junker, Y. P. Chen, T. A. Corcovilos, and R. G. Hulet, Phys. Rev. Lett. 102, 090402 (2009).
[2] Dan M. Stamper-Kurn and Masahito Ueda, Rev. Mod. Phys. 85, 1191 (2013).
[3] J. Stenger, S. Inouye, D. M. Stamper-Kurn, H.-J. Miesner, A. P. Chikkatur and W. Ketterle, Nature 396, 345 (1998).
[4] Simon Murmann, Andrea Bergschneider, Vincent M.
[5] Cheng Chin, Rudolf Grimm, Paul Julienne, and Eite Tiesinga, Rev. Mod. Phys. 82, 1225 (2010).

[6] P. O. Fedichev, Yu. Kagan, G. V. Shlyapnikov, and J. T. M. Walraven, Influence of nearly resonant light on the scattering length in low-temperature atomic gases, Phys. Rev. Lett. 77, 2913 (1996).

[7] L.-M. Duan, E. Demler, M. D. Lukin, Phys. Rev. Lett. 91, 090402 (2003); [arXiv:cond-mat/0210564]

[8] R. Ciuryłło, E. Tiesinga, and P. S. Julienne, Phys. Rev. A 71, 030701(R) (2005).

[9] Pascal Naidon and Françoise Masnou-Seeuws, Phys. Rev. A 73, 043611 (2006).

[10] Thorsten Köhler, Krzysztof Góral, and Paul S. Julienne, Rev. Mod. Phys. 78, 1311 (2006).

[11] M. Theis, G. Thalhammer, K. Winkler, M. Hellwig, G. Ruff, R. Grimm, J. Hecker Denschlag, Tuning the scattering length with an optically induced Feshbach resonance, Phys. Rev. Lett. 93, 123001 (2004); [arXiv:cond-mat/0404514]

[12] John L. Bohn and P. S. Julienne, Prospects for influencing scattering lengths with far-off-resonant light, Phys. Rev. A 56, 1486 (1997).

[13] Felix H.J. Hall, Mireille Aymar, Nadia Boulonfa-Maafa, Olivier Dulieu, Stefan Willitsch, Light-assisted ion-neutral reactive processes in the cold regime: radiative molecule formation vs. charge exchange, Phys. Rev. Lett. 107, 243202 (2011); [arXiv:1108.3739]

[14] Adam Micah Kaufman, Radiofrequency dressing of atomic Feshbach resonances, Submitted to the Department of Physics of Amherst College in partial fulfilment of the requirements for the degree of Bachelors of Arts, 2009.

[15] S. Blatt, T. L. Nicholson, B. J. Bloom, J. R. Williams, J. W. Thomsen, P. S. Julienne, and J. Ye, Phys. Rev. Lett. 107, 073202 (2011).

[16] A. C. Hewson, The Kondo Problem to Heavy Fermions (Cambridge University Press, Cambridge, 1993).

[17] J. Bauer, C. Salomon and E. Demler, Phys. Rev. Lett. 111, 215304 (2013).

[18] I. Kuzmenko, T. Kuzmenko, Y. Avishai and K. A. Kikoin, Phys. Rev. B 91, 165131 (2015); [arXiv:1402.0187]

[19] Igor Kuzmenko, Tetyana Kuzmenko, Yshai Avishai and Gyu-Boong Jo, Phys. Rev. B 93, 115143 (2016); [arXiv:1512.00978]

[20] Igor Kuzmenko, Tetyana Kuzmenko, Yshai Avishai and Gyu-Boong Jo, to be published in Phys. Rev. B.

[21] Ren Zhang, Deping Zhang, Yanting Cheng, Wei Chen, Peng Zhang, and Hui Zhai, Phys. Rev. A 93, 043601 (2016); [arXiv:1509.01350]

[22] Yanting Cheng, Ren Zhang, Peng Zhang, and Hui Zhai, Phys. Rev. A 96, 063605 (2017); [arXiv:1705.06878]

[23] Luis Rieger, Nelson Darkwah Oppong, Moritz Hfer, Diogo Rio Fernandes, Immanuel Bloch and Simon Flling, [arXiv:1708.03810]

[24] Qing Ji, Ren Zhang, Xiang Zhang, Wei Zhang, [arXiv:1809.00471]

[25] W. F. Meggers and J. L. Tech, J. Res. Natl. Bur. Stand. (U.S.) 83, 13 (1978).

[26] T. Andersen, “Atomic negative ions: Structure, dynamics and collisions”. Physics Reports 394, 157 (2004).

[27] Masaaki Kitagawa, Katsunari Enomoto, Kentaro Kasa, Yoshiro Takahashi, Roman Ciuryłło, Pascal Naidon, and Paul S. Julienne, Phys. Rev. A 77, 012719 (2008); [arXiv:0708.0752]

[28] G. F. Gribakin and V. V. Flambaum, Phys. Rev. A 48, 546 (1993).

[29] F. Scazza, C. Hofrichter, M. Fofer, P. C. De Groot, I. Bloch, and S. Folling, Nature Physics 10, 779 (2014); [ibid: correction notice, (2015)].

[30] Comparison between the CIR and optical techniques will be assessed elsewhere.

[31] From the present analysis we now conclude that the claim that antiferromagnetic exchange exists also in the absence of laser field as reported in Ref. [19] should be retracted. Fortunately, it can be cured following the analysis detailed in the present note. A proper corrigendum will shortly be reported.
About spontaneous emission

Spontaneous emission is dangerous since it heats the atomic gas. In order to avoid this, we need to organize the procedure as follows:

1. The frequency of detuning from the resonant frequency is needed to be large enough so that absorption/emission of the photon is forbidden by energy conservation.

2. The $^1S_0$ and $^3P_0$ quantum states have different electronic principal quantum numbers and spin states. Electric dipole transition between the singlet states $^1S_0$ and $^3P_0$ is forbidden. Magnetic dipole transition can change the spin, but not the principal quantum number. Therefore, spontaneous quantum transition from $^3P_0$ to $^1S_0$ state of Yb is virtually forbidden (or at least, the lifetime of the $^3P_0$ state is very long).

Note that the same problem exists also for a "bare" mixture of Yb($^1S_0$) and Yb($^3P_0$) atoms (without an additional laser light). However, applying an additional laser radiation makes the exchange interaction stronger and (possible) the lifetime shorter. So far, we do not know how to calculate the lifetime.

In an experimental work [arXiv:1708.03810] the authors write: “For the scattering of the ionization beam, we will try to estimate the scattering rates and compare it with our trap depths. Probably, we will have to choose a significant detuning as not to heat the atoms out of the trapping confinement.”