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Implementation of quantum search scheme by adiabatic passage in a cavity–laser–atom system*

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This paper proposes a scheme for implementing the adiabatic quantum search algorithm of different marked items in an unsorted list of N items with atoms in a cavity driven by lasers. N identical three-level atoms are trapped in a single-mode cavity. Each atom is driven by a set of three pulsed laser fields. In each atom, the same level represents a database entry. Two of the atoms are marked differently. The marked atom has an energy gap between its two ground states. The two different marked states can be sought out respectively starting from an initial entangled state by controlling the ratio of three pulse amplitudes. Moreover, the mechanism, based on adiabatic passage, constitutes a decoherence-free method in the sense that spontaneous emission and cavity damping are avoided since the dynamics follows the dark state. Furthermore, this paper extends the algorithm with m ($m > 2$) atoms marked in an ideal situation. Any different marked state can be sought out.

Keywords: adiabatic quantum search algorithm, cavity–laser–atom system, marked state
PACC: 0365, 4250

1. Introduction

Grover showed how a quantum computer can find a single marked object in a database of size N by using only $O\left(\sqrt{N}\right)$ queries of the oracle that identifies the object.[1] His result was generalized to the case of finding one object in a subset of marked elements. Classically, this problem requires approximately N operations. Even the most efficient algorithm still needs on average $\frac{N}{2}$ operations. Searching an unsorted database has always been a very important mathematical problem. Moreover, many techniques based on the search algorithm are widely used in our life, and it can have utility for a wide range of problems,[2] so the quantum search algorithm becomes more and more important. Much effort has been devoted to the achievement of these quantum algorithms theoretically and experimentally by using NMR systems,[3−6] trapped ions,[7,8] optics,[9−12] cavity QED[13−15] and superconducting quantum interference devices in a cavity.[16]

Quantum adiabatic algorithms, first conceived in Ref.[17], have been developed as an alternative to the traditional circuit model of quantum computation. An important aspect of the adiabatic approach is its expected resilience to various kinds of open system effects, which makes it a promising candidate for robust quantum computation. Adiabatic versions of the time-continuous Grover algorithm have been proposed[17−19] mainly to take advantage of the robustness of adiabatic passage with respect to fluctuations of the external control fields as well as the imperfect knowledge of the model. The robustness of the adiabatic quantum search to decoherence in the instantaneous eigenbasis of the search Hamiltonian was examined.[20,21]

The implementation of the adiabatic Grover algorithm based on a physical system has been proposed by Daems and Guérin.[22] It is formulated with a Hamiltonian $H = H_0 + V(t)$, where H_0 is given and considered as an oracle, while $V(t)$ is slowly varying in time such that there is an instantaneous eigenvector adiabatically connecting the initial superposition $|w\rangle$ to the marked state $|m\rangle$. To our knowledge, this is the first proposed implementation of the adiabatic Grover algorithm. However, there is only one marked atom and the other $N-1$ atoms are the same. In practice, there will be many different items in an unsorted database, and each of them needs to be sought out.

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In this paper, motivated by the previous works, we discuss how to deterministically search out any different single marked state in the cavity–laser–atom system with $m\,(m \geq 2)$ atoms marked based on a previous idea.[22] We first marked two atoms and let other $N-2$ atoms be the same. Our task is to find any different single marked item among N items stored in an unsorted database. Then we extend the system to mark m $(m > 2)$ atoms. We show that any of the m different marked items can be sought out.

2. The system

The involved atomic levels and transitions are depicted in Fig.1. N identical $\wedge$-type atoms are trapped in a single-mode cavity. Each atom has one excited state $|e\rangle$ and two ground states $|g\rangle$ and $\left|g^{/}\right\rangle$. The $|e\rangle \leftrightarrow |g\rangle$ transition is coupled resonantly to the single-mode cavity with Rabi frequency G. The transitions $|e\rangle \leftrightarrow \left|g^{/}\right\rangle$ of two marked atoms are driven resonantly through two lasers with Rabi frequencies $\Omega^{/}$ and $\Omega^{//}$, respectively. The laser of Rabi frequency Ω is resonant with the $|e\rangle \leftrightarrow \left|g^{/}\right\rangle$ transition for the unmarked atoms. The two ground states $|g\rangle$ and $\left|g^{/}\right\rangle$ of the marked (unmarked) atoms are non-degenerate (degenerate). The states $\left|g^{/}\right\rangle$ of two marked atoms are shifted in energy by δ_1 and δ_2 respectively with respect to those of the unmarked atoms. We assume that the energy of the states $\left|g^{/}\right\rangle$ and $|g\rangle$ of the unmarked atoms is zero.

This can be realized in practice by considering Zeeman states and laser and cavity fields of appropriate polarizations. The polarizations and the frequencies are such that each field drives an unique transition by a one-photon resonant process. As a consequence, the Stark shifts, which would add detrimental phases, can be neglected here. The states $\left|g^{/}\right\rangle$ of the N atoms represent the database entries. The states $|g\rangle$ allow for the coupling of all the atoms through the exchange of a single photon with the cavity (see Fig.2(a)). We shall start from the initial state $\left|g^{/}, 0\right\rangle \equiv \frac{1}{\sqrt{N}} \sum_{j=1}^{N} \left|g_j^{/}, 0\right\rangle$ which features a collective superposition of three types of atomic ground states. Each atom is denoted as a subscript $1, \ldots, N$. We choose tags N and $N-1$ to label the two marked atoms respectively.

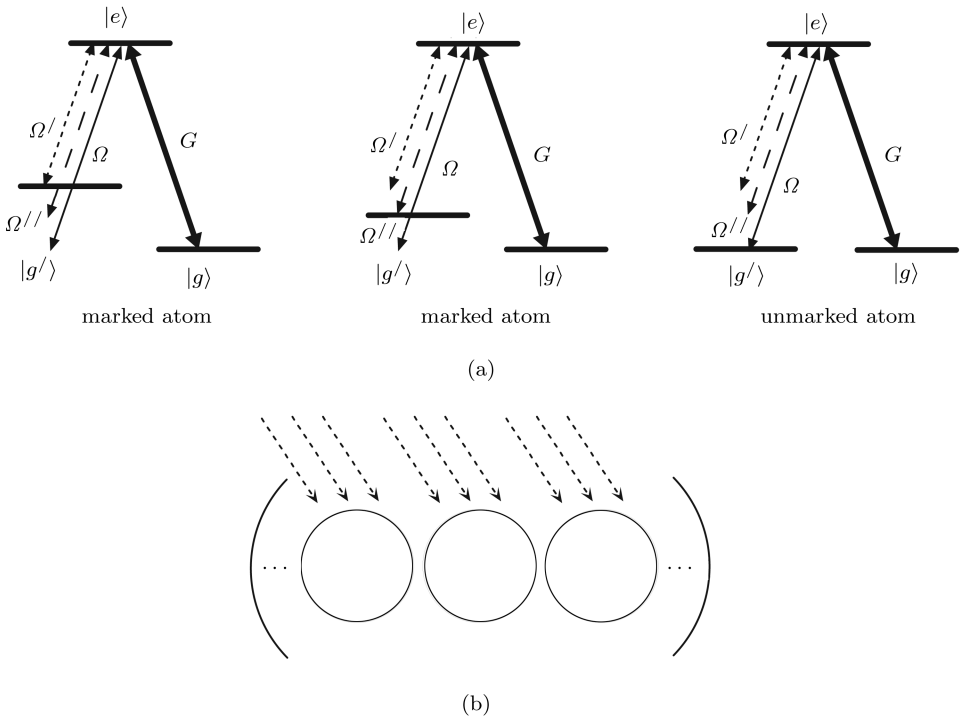


Fig.1. (a) Linkage pattern for the individual atoms. The cavity G, laser Ω, laser $\Omega^{/}$ and laser $\Omega^{//}$ Rabi frequencies are featured by thick, thin, dotted, and dashed arrows, respectively. (b) Representation of the atoms register trapped in a single-mode optical cavity. The atoms are represented by circles, laser fields by arrows.

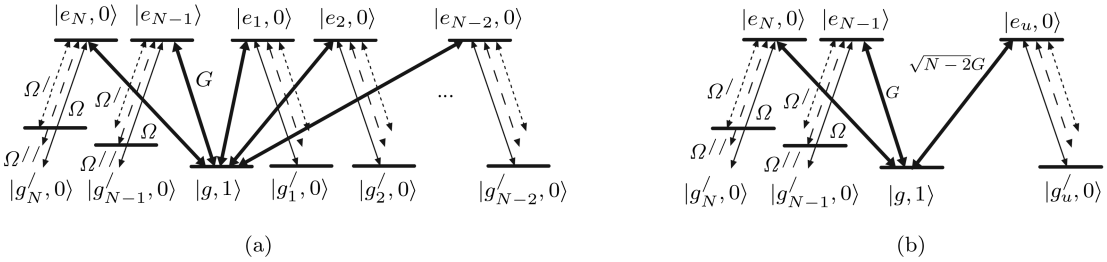


Fig.2. (a) Coupling scheme in the cavity. (b) Equivalent scheme. The effective cavity Rabi frequency to the collective unmarked excited state is $\sqrt{N-2}G$.

3. Implementation of the adiabatic quantum search scheme

We consider a cavity mode of frequency w and coupling strength G together with three lasers of frequencies $w, w-\delta_1,\ w-\delta_2$ and pulse shapes $\Omega(t)$, $\Omega^{/}(t)$, $\Omega^{//}(t)$ which do not grow with N. Spontaneous emission can be neglected when the condition $(\Omega_{\max}T)^2 \gg \Gamma T$ is fulfilled.

For this case, the Hamiltonian of the whole system can be written as (assuming $\hbar = 1$)

$$H = H_0 + H_{\mathrm{I}}, \tag{1}$$

where

$$H_0 = \delta_1 \left|g_N^{/}\right\rangle \left\langle g_N^{/}\right| + \delta_2 \left|g_{N-1}^{/}\right\rangle \left\langle g_{N-1}^{/}\right| + w\sum_{j=1}^{N} |e_j\rangle\langle e_j| + wa^{+}a, \tag{2}$$

$$H_{\mathrm{I}} = Ga\sum_{j=1}^{N} |e_j\rangle\langle g_j| + [\Omega \mathrm{e}^{\mathrm{i}wt} + \Omega^{/}\mathrm{e}^{\mathrm{i}(w-\delta_1)t} + \Omega^{//}\mathrm{e}^{\mathrm{i}(w-\delta_2)t}] \times \sum_{j=1}^{N} \left|g_j^{/}\right\rangle\langle e_j| + \mathrm{H.C.}, \tag{3}$$

where H_0 is the energy of the system consisting of N atoms and the single-mode cavity, H_{I} is the interaction of atoms with cavity and lasers. H.C. denotes the Hermitian conjugate. a^{+} and a are creation and annihilation operators associated with the single-mode cavity respectively. The energy of the excited state $|e\rangle$ is w (the energy reference is taken for the ground states $\left|g^{/}\right\rangle$ and $|g\rangle$ of the unmarked atoms). We shall focus on the system associated with a single photon in the cavity. When the N atoms are in their ground state $|g\rangle$, the corresponding multipartite state $|g_1 \ldots g_N\rangle \otimes |1\rangle$ is denoted by $|g,1\rangle$. The excited state $|e\rangle$ of any of the N atoms, say atom j, can be reached by absorption of the cavity photon while the other atoms remain in their ground state $|g\rangle$; the corresponding multipartite state is $|e_j,0\rangle \equiv |g_1 \ldots g_{j-1} e_j g_{j+1} \ldots\rangle \otimes |0\rangle$. The state $|e_j,0\rangle$ can also be reached by absorption of one laser photon when any of the atoms, say atom j, is in the ground state $\left|g^{/}\right\rangle$ while the other atoms are in the ground state $|g\rangle$: $\left|g_j^{/},0\right\rangle \equiv \left|g_1 \ldots g_{j-1} g_j^{/} g_{j+1} \ldots g_N\right\rangle \otimes |0\rangle$.

The Hamiltonian in the interaction picture

$$V(t) = U_0^{\dagger}(t) H_{\mathrm{I}} U_0(t), \tag{4}$$

with

$$U_0(t) = \mathrm{e}^{-\mathrm{i}H_0 t} = \mathrm{e}^{-\mathrm{i}t\left(\delta_1 \left|g_N^{/}\right\rangle\left\langle g_N^{/}\right| + \delta_2 \left|g_{N-1}^{/}\right\rangle\left\langle g_{N-1}^{/}\right| + w\sum\limits_{j=1}^{N} |e_j\rangle\langle e_j| + wa^{+}a\right)}, \tag{5}$$

reads

$$V = Ga\sum_{j=1}^{N}|e_j\rangle\langle g_j| + (\Omega + \Omega^{/}\mathrm{e}^{-\mathrm{i}\delta_1 t} + \Omega^{//}\mathrm{e}^{-i\delta_2 t})\sum_{j=1}^{N-2}\left|g_j^{/}\right\rangle\langle e_j| + [\Omega\mathrm{e}^{\mathrm{i}\delta_1 t} + \Omega^{/} + \Omega^{//}\mathrm{e}^{\mathrm{i}(\delta_1-\delta_2)t}]\left|g_N^{/}\right\rangle\langle e_N|$$
$$+[\Omega\mathrm{e}^{\mathrm{i}\delta_2 t} + \Omega^{/}\mathrm{e}^{\mathrm{i}(\delta_2-\delta_1)t} + \Omega^{//}]\left|g_{N-1}^{/}\right\rangle\langle e_{N-1}| + \mathrm{H.C.}. \quad (6)$$

Among the unmarked atoms, none should play a privileged role. Hence we shall consider them collectively and label the corresponding state with a subscript u. The states $\left|g_j^{/},0\right\rangle$ ($|e_j,0\rangle$) of the unmarked atoms form the collective unmarked ground state $\left|g_u^{/},0\right\rangle$ (excited state $|e_u,0\rangle$), i.e., the entangled states $\left|g_u^{/},0\right\rangle = \frac{1}{\sqrt{N-2}}\sum_{j=1}^{N-2}\left|g_j^{/},0\right\rangle$ and $|e_u,0\rangle = \frac{1}{\sqrt{N-2}}\sum_{j=1}^{N-2}|e_j,0\rangle$.

In this new basis, the Hamiltonian V restricted to the subspace spanned by the states $\left|g_u^{/},0\right\rangle$, $\left|g_N^{/},0\right\rangle$, $\left|g_{N-1}^{/},0\right\rangle$, $|g,1\rangle$, $|e_u,0\rangle$, $|e_N,0\rangle$, $|e_{N-1},0\rangle$ reads

$$V = \begin{pmatrix} 0 & 0 & 0 & 0 & \Sigma & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & \Sigma^{/} & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 & \Sigma^{//} \\ 0 & 0 & 0 & 0 & \sqrt{N-2}G & G & G \\ \Sigma^{*} & 0 & 0 & \sqrt{N-2}G & 0 & 0 & 0 \\ 0 & \Sigma^{/*} & 0 & G & 0 & 0 & 0 \\ 0 & 0 & \Sigma^{//*} & G & 0 & 0 & 0 \end{pmatrix}, \quad (7)$$

with $\Sigma = \Omega + \mathrm{e}^{-\mathrm{i}\delta_1 t}\Omega^{/} + \mathrm{e}^{-\mathrm{i}\delta_2 t}\Omega^{//}$, $\Sigma^{/} = \mathrm{e}^{\mathrm{i}\delta_1 t}\Omega + \Omega^{/} + \mathrm{e}^{\mathrm{i}(\delta_1-\delta_2)t}\Omega^{//}$, and $\Sigma^{//} = \mathrm{e}^{\mathrm{i}\delta_2 t}\Omega + \mathrm{e}^{\mathrm{i}(\delta_2-\delta_1)t}\Omega^{/} + \Omega^{//}$. This Hamiltonian is represented in Fig.2(b). The non-resonant components in Σ, $\Sigma^{/}$ and $\Sigma^{//}$ vary very rapidly and their average over a time scale $|\delta_1 - \delta_2|^{-1}$ is zero. Hence we can get the resonant approximation: $\bar{\Sigma} \approx \Omega$, $\bar{\Sigma}^{/} \approx \Omega^{/}$ and $\bar{\Sigma}^{//} \approx \Omega^{//}$.

Our aim is to change the population by driving adiabatically the initial entangled state $\left|g^{/},0\right\rangle$, which corresponds to the superposition of an unmarked collective state $\left|g_u^{/}\right\rangle$ and two marked states $\left|g_N^{/},0\right\rangle$ and $\left|g_{N-1}^{/},0\right\rangle$, to the single marked state $\left|g_N^{/},0\right\rangle$ or $\left|g_{N-1}^{/},0\right\rangle$ (see Fig.2(b)). The state $\left|g^{/},0\right\rangle$ gives no privileged role to any of the N states $\left|g_j^{/},0\right\rangle$. The population transfer mechanism is most easily revealed in the basis of the instantaneous eigenstates of $V(t)$. The Hamiltonian matrix of $V(t)$ has a dark state (the eigenstate of the Hamiltonian with zero energy eigenvalue):

$$D_1(t) = -\frac{\sqrt{N-2}G\Omega^{/}\Omega^{//}}{\sqrt{(N-2)G^2\Omega^{/2}\Omega^{//2} + G^2\Omega^2\Omega^{//2} + G^2\Omega^2\Omega^{/2} + \Omega^2\Omega^{/2}\Omega^{//2}}}\left|g_u^{/},0\right\rangle$$
$$-\frac{G\Omega\Omega^{//}}{\sqrt{(N-2)G^2\Omega^{/2}\Omega^{//2} + G^2\Omega^2\Omega^{//2} + G^2\Omega^2\Omega^{/2} + \Omega^2\Omega^{/2}\Omega^{//2}}}\left|g_N^{/},0\right\rangle$$
$$-\frac{G\Omega\Omega^{/}}{\sqrt{(N-2)G^2\Omega^{/2}\Omega^{//2} + G^2\Omega^2\Omega^{//2} + G^2\Omega^2\Omega^{/2} + \Omega^2\Omega^{/2}\Omega^{//2}}}\left|g_{N-1}^{/},0\right\rangle$$
$$+\frac{\Omega\Omega^{/}\Omega^{//}}{\sqrt{(N-2)G^2\Omega^{/2}\Omega^{//2} + G^2\Omega^2\Omega^{//2} + G^2\Omega^2G^{/2} + \Omega^2\Omega^{/2}\Omega^{//2}}}|g,1\rangle\,. \quad (8)$$

The dynamics will follow the atomic dark state if the Hamiltonian varies sufficiently slowly in time. Hence, the excited atomic state is never populated (in the adiabatic limit). Moreover, the projections of the dark states into the excited cavity photon states can be made negligible if $G \gg \Omega_{\max}$.[23−25] In this case, the proposed mechanism is a decoherence-free method in the sense that the process is not sensitive to spontaneous emission from the atomic excited states as well as the lifetime of photons in the optical cavity.

Now let us describe the process of the adiabatic quantum search algorithm of the different single marked states in detail. Firstly, we switch on Ω, $\Omega^{/}$ and $\Omega^{//}$ simultaneously, so the instantaneous eigenstate $|D_1(t)\rangle$ is initially in the state $\left|g^{/},0\right\rangle$. Then we switch off $\Omega^{/}(\Omega^{//})$ before Ω and $\Omega^{//}(\Omega^{/})$. Finally, the instantaneous eigenstate $\|D_1(t)\rangle$ coincides with the marked state $|g_N^{/},0\rangle(|g_{N-1}^{/},0\rangle)$. Thus, we can seek out the two different marked states respectively. The characteristic time to achieve such a process was shown to grow as $\sqrt{N}$ in order to satisfy adiabaticity.[22]

4. Extension to the system with m $(m > 2)$ atoms marked

In this section, we show that the m different marked states can be deterministically searched out respectively in an ideal situation. m atoms are marked differently in the single-mode cavity. The energy of the state $\left|g^{/}\right\rangle$ of the mth marked atom is shifted by an amount δ_m with respect to that of the unmarked atoms. The energy of the state $\left|g^{/}\right\rangle$ of unmarked atoms is set to zero. Each atom (labelled by j) is resonantly driven by a set of $(m+1)$ pulsed laser fields Ω_0, Ω_1, $\Omega_2 \ldots \Omega_m$ and by the cavity mode G. In the interaction picture, the Hamiltonian for system of cavity–laser–atom can be described by

$$\begin{aligned} V_2 = {} & Ga\sum_{j=1}^{N}|e_j\rangle\langle g_j| + (\Omega_0 + \Omega_1 e^{-i\delta_1 t} + \Omega_2 e^{-i\delta_2 t} + \ldots + \Omega_m e^{-i\delta_m t}) \sum_{j=1}^{N-m}\left|g_j^{/}\right\rangle\langle e_j| \\ & + [\Omega_0 e^{i\delta_1 t} + \Omega_1 + \Omega_2 e^{i(\delta_1-\delta_2)t} + \ldots + \Omega_m e^{i(\delta_1-\delta_m)t}]\left|g_N^{/}\right\rangle\langle e_N| \\ & + [\Omega_0 e^{i\delta_2 t} + \Omega_1 e^{i(\delta_2-\delta_1)t} + \Omega_2 + \ldots + \Omega_m e^{i(\delta_2-\delta_m)t}]\left|g_{N-1}^{/}\right\rangle\langle e_{N-1}| \\ & + \ldots + [\Omega_0 e^{i\delta_m t} + \Omega_1 e^{i(\delta_m-\delta_1)t} + \Omega_2 e^{i(\delta_m-\delta_2)t} + \ldots + \Omega_m]\left|g_{N-m+1}^{/}\right\rangle\langle e_{N-m+1}| + \text{H.C.}, \end{aligned} \tag{9}$$

where $m < N$. Subscripts j and m represent the jth and mth atom, respectively. The subspace can be written as follows: $\left|g_u^{/},0\right\rangle$, $\left|g_N^{/},0\right\rangle$, $\left|g_{N-1}^{/},0\right\rangle \ldots \left|g_{N-m+1}^{/}\right\rangle$, $|g,1\rangle$, $|e_u,0\rangle$, $|e_N,0\rangle$, $|e_{N-1},0\rangle \ldots |e_{N-m+1}\rangle$, where $\left|g_u^{/},0\right\rangle = \frac{1}{\sqrt{N-m}}\sum_{j=1}^{N-m}\left|g_j^{/},0\right\rangle$, $|e_u,0\rangle = \frac{1}{\sqrt{N-m}}\sum_{j=1}^{N-m}|e_j,0\rangle$. Here collective state $\left|g_u^{/},0\right\rangle$ interacts with collective state $|e_u,0\rangle$ with an effective coupling strength $\sqrt{N-m}G$. The transition $|g_{N-m+1}^{/}\rangle \leftrightarrow |e_{N-m+1}\rangle$ of the marked atom is driven resonantly through the mth laser with Rabi frequency Ω_m. The Hamiltonian V_2 can be written as a matrix in terms of the subspace. Corresponding to the subspace, the dark state (the eigenstate of the Hamiltonian V_2 with zero energy eigenvalue) can be written as

$$\begin{aligned} D_2(t) = {} & -(\sqrt{N-m}G\Omega_1\Omega_2\ldots\Omega_m C)\left|g_u^{/},0\right\rangle - (G\Omega_0\Omega_2\ldots\Omega_m C)\left|g_N^{/},0\right\rangle \\ & - \ldots - (G\Omega_0\Omega_1\ldots\Omega_{m-1}C)\left|g_{N-m+1}^{/},0\right\rangle + (\Omega_0\Omega_1\ldots\Omega_m C)\left|g,1\right\rangle. \end{aligned} \tag{10}$$

Where $C = \dfrac{1}{\sqrt{(N-m)G^2\Omega_1^2\Omega_2^2\ldots\Omega_m^2 + G^2\Omega_0^2\Omega_2^2\ldots\Omega_m^2 + \ldots + G^2\Omega_0^2\Omega_1^2\ldots\Omega_{m-1}^2 + \Omega_0^2\Omega_1^2\ldots\Omega_m^2}}$.

The condition $G \gg \Omega_{\max}$ has to be satisfied such that the cavity mode is negligibly populated during the interaction with the pulses. Requiring the instantaneous dark state $|D_2(t)\rangle$ to coincide at the initial time with the uniform superposition of all ground states of the subspace, we find that $(m+1)$ pulses must be switched on simultaneously, and to find out the mth marked state at the final time, the pulse Ω_m is to be turned off before other pulses. Then the mth marked state can be searched out and the search scheme can be realized.

5. Discussion and conclusion

Finally, we will briefly discuss the experimental feasibility of our scheme. We can trap atoms in a cavity by using a standing-wave dipole-force trap for the experimental implementation of the robust processes.[26] As a realistic atomic level scheme, we can consider the $2^3S_1 - 2^3P_0$ transition in metastable helium, of linewidth $\Gamma = 10^7 \mathrm{s}^{-1}$ (see for instance Ref.[27]). The first excited state of helium (2^3S_1, some 20 eV above the ground state) has a very long (metastable) lifetime of $\tau \sim 8000\,\mathrm{s}$.[28] Thus 2^3S_1 metastable helium can effectively be employed as a 'ground-state' species. Energy shifts of the magnetic sublevels of $^3S_1$ helium caused by optical pumping have been well discussed in Ref.[29]. Taking into account the losses of the cavity (characterized by the decay rate κ of the cavity field) and of the excited states, we have to satisfy the adiabatic conditions $\Omega_{\max}T$, $GT \gg 1$ and $(\Omega_{\max}T)^2 \gg \Gamma T$, $(GT)^2 \gg \kappa T$. The latter is satisfied for $\Omega_{\max} \gg \Gamma$ and $G \gg \kappa$. For a typical pulse duration of $T = 50\,\mathrm{ns}$, we use $\Omega_{\max} = 1.2 \times 10^8 \mathrm{s}^{-1}$, which is achievable experimentally (see for instance Ref.[30]). We use a cavity coupling $G = 5.2 \times 10^8 \mathrm{s}^{-1}$, more than four times larger than $\Omega_{\max}$ to have a small population in the cavity field. Such a strong coupling has been achieved in experiments with atoms in an optical cavity trapped with a duration of the order of one second.[31,32]

In conclusion, we have proposed a scheme to implement the quantum search algorithm of any different single marked item in an unsorted list of N items by the adiabatic passage of dark state in a cavity-laser-atom system. We have shown that the marked atoms can be exploited to realize scalability. The proposed scheme offers a promising approach to the realization of large-scale adiabatic quantum search. As predicted, the transfer to the marked state is very efficient. In our scheme, there are some advantages as follows. (i) The collective excited states $\left|e_N^{\prime}, 0\right\rangle$, $\left|e_{N-1}^{\prime}, 0\right\rangle$ and $\left|e_u^{\prime}, 0\right\rangle$ would not be excited, and the operation only involves three ground states $\left|g_u^{\prime}, 0\right\rangle$, $\left|g_N^{\prime}, 0\right\rangle$ and $\left|g_{N-1}^{\prime}, 0\right\rangle$. Therefore, we can omit approximately the effects of decoherence due to the spontaneous emission of the excited states. (ii) The amplitude of the laser field and interaction time need not be accurately controlled, as long as these parameters can satisfy the adiabatic condition. Thus, the operation is insensitive to certain errors, such as the amplitude fluctuations of the laser fields. (iii) Different from Ref.[22], we realize adiabatic quantum search of any different single marked item in a database of unsorted N items. In addition, in the scheme of marking two atoms, three laser fields are applied to each atom. In order to search any different marked state, we only need to control the three lasers. Furthermore, we have theoretically discussed an extension of the method to the system with m ($m > 2$) atoms marked. The operation process will not be changed with the increasing number of lasers. The property can significantly reduce the number of operations for searching the mth different marked state in an experiment. Therefore, we believe that it is feasible with advanced techniques to realize the quantum processes in our paper.

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