Coherently manipulating flying qubits in a quantum wire with a magnetic impurity

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We study the effect of a magnetic impurity with spin-half on a single propagating electron in a one-dimensional model system via the tight-binding approach. Due to the spin-dependent interaction, the scattering channel for the flying qubit is split, and its transmission spectrum is obtained. It is found that, the spin orientation of the impurity plays the role as a spin state filter for a flying qubit.

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I. INTRODUCTION

Qubits are the quantum state of given physical systems. There are two kinds of qubits, the first kind is called stationary qubits which are fixed into space, the second kind is the so-called flying qubits whose position changes in time. To faithfully transfer quantum information, individual qubit control would be desirable. Controlling qubits are implemented by a sequence of logical operations or quantum gates. For stationary qubits, quantum gates are performed by choosing a proper time, however, a logic operation has to be fixed into space for flying qubits.

Nowadays, more attention has been paid on the use of the flying qubit due to the advantage that it allows one to entangle distant stationary qubits that never interacted directly.\cite{1,2}. Therefore, there are with keen interest in seeking new device for controlling flying qubits.\cite{3,4,5,6}. Photon is the archetypal flying qubit, which is special ideal for long distance communication due to its high speed, strong stability and minor loss. Therefore quantum devices that enable new filtering and switching functions have been proposed, such as single photon transistor\cite{7,8}, quantum switch\cite{4,5,6}. The spin of a propagating electron is an alternative leading candidate for a flying qubit\cite{12,13} in short distance communication mostly due to the intrinsic nature of electron spin, namely that the spin degree of freedom is well isolated from the environment. The ability to manipulate an mobile spin qubits is indispensable for the potential application in quantum information processing. A singlet spin filter has been proposed by the Coulomb interaction between a flying qubit and the trapped electron in a weak confining potential\cite{14}. Later on it was generalized to a potential well with multiple bound orbitals\cite{15}. In this paper, we propose a scheme whereby a single propagating electron is subject to a spin-spin interaction with a magnetic impurity in a quantum wire. The local external magnetic field removes the degeneracy. Then the spin-dependent scattering behavior of the propagating electron has been investigated via the tight-binding approach. Due to the spin-flip interaction between a single propagating electron and the magnetic impurity, a total reflection of the propagating electron has been found by a proper choice of the injection energy of the electron.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig1}
\caption{(Color online). Schematic of the perfect quantum wire in one dimension, where a magnetic impurity with spin-1/2 is embedded in the 0th site.}
\end{figure}

\begin{equation}
H_c = \sum_{\sigma} H_{qw}^{\sigma} + H_{fr}^{\sigma} + H_{ex},
\end{equation}

\section{A FLYING QUBIT WITH A MAGNETIC IMPURITY}

The model we considered in this paper is shown in Fig.\textsuperscript{1}. Thanks to the success of nanofabrication techniques in producing extremely small quantum objects, this kind of quantum wires are now close to reality either by small semiconductor and metal structures as a fabricated quantum dot array\cite{16,17}, or by the scanning tunnel microscope\cite{18,19} as the atomic chain.

The system in Fig.\textsuperscript{1} is a one-dimensional wire, which also has an impurity embedded in a host lattice. The Hamiltonian of this system consists three parts
where

\[ H_{qw}^{e\sigma} = \sum_j \omega_j a_{j\sigma}^\dagger a_{j\sigma} + \sum_j \xi \left( a_{j\sigma}^\dagger a_{j+1\sigma} + \text{h.c.} \right), \]  
\[ H_{fr}^{\sigma} = \sum_{\sigma} (\omega_0 - \omega) a_{0\sigma}^\dagger a_{0\sigma} + \Delta \left( a_{0e}^\dagger a_{0e} - a_{0g}^\dagger a_{0g} \right) + 2\Omega_l S_I^z, \]
\[ H_{ex} = -2J_z \left( a_{0e}^\dagger a_{0e} - a_{0g}^\dagger a_{0g} \right) S_I^z - J \left( a_{0e}^\dagger a_{0g} S_I^- + \text{h.c.} \right). \]

Here the superscript \( \sigma \in \{ e, g \} \) in \( H_{qw}^{e\sigma} \) describes the spin degree of freedom, where \( e \) (\( g \)) denotes spin up (down). The host lattice is modeled as tight-binding Hamiltonian, which is represented by a sequence of potential sites with on-site energies \( \omega \) except the one at \( j = 0 \), and constant amplitude \( \xi \) of hopping between the neighboring sites. \( a_{j\sigma}^\dagger \) (\( a_{j\sigma} \)) is the creation (annihilation) operator for the 1D wire. The magnetic impurity of spin \( S_I \) is located only at one site, which is chosen as the origin. A given spin state of the impurity is changed by the spin rising operator \( S_I^+ \) and lowering operator \( S_I^- \). In Eq. (2), Hamiltonian \( H_{fr}^{\sigma} \) has taken the interaction energy of the spin magnetic moments with the magnetic field into account through the terms \( \Delta \left( a_{0e}^\dagger a_{0e} - a_{0g}^\dagger a_{0g} \right) \) for the propagating electrons and \( 2\Omega_l S_I^z \) for the impurity. Here, \( \Delta \) and \( \Omega_l \) are the Zeeman splitting of the energy levels in the presence of an external magnetic field along the \( z \)-direction. The exchange interaction among spins, which is a short range point-like interaction localized at the point \( j = 0 \), is described by Hamiltonian \( H_{ex} \). Here we introduce the anisotropic spin-spin interaction with XXZ-type. Since we are interested only in the coupling to single impurity, the interaction here is written in terms of the single spin operator above, rather than the second-quantized description of the impurity. By the Fourier transform, the Hamiltonian for the conduction electrons reads

\[ H_{qw}^{e\sigma} = \sum_k (\omega + 2\xi \cos k) a_{k\sigma}^\dagger a_{k\sigma} \]

where \( a_{k\sigma}^\dagger \) and \( a_{k\sigma} \) are the usual creation and annihilation operators of the conduction electron with wave vectors \( k \) and spin \( \sigma \). And the anisotropic spin-spin interaction reads

\[ H_{ex} = -2J_z \sum_{kk'} \left( a_{k\sigma}^\dagger a_{k'\sigma}^\dagger a_{k'\sigma} - a_{k\sigma}^\dagger a_{k\sigma} \right) S_I^z \]

\[ -J \sum_{kk'} \left( a_{k\sigma}^\dagger a_{k'\sigma}^\dagger S_I^- + \text{h.c.} \right) \]

where \( N \) is the number of sites in the lattice. \( \sum_{\sigma} H_{qw}^{e\sigma} + H_{ex} \) is the Hamiltonian for the system consisting of a localized spin in interaction with the conduction band, which is a simpler Kondo model when \( J_z = J \).

III. CONTROLLING THE PROPAGATING ELECTRON IN 1D QUANTUM WIRE

A flying qubit is a physical realization of a qubit which moves freely and allows information to be transported from one location to another. Obviously, the spin of an electron propagating along the host 1D structure acts as a flying qubit, it also interacts and exchanges information with the magnetic impurity due to the exchange interaction. For the sake of simplicity, we assume the magnetic impurity is a spin-1/2 particle in the following discussion. The single-electron particle picture is employed to study the quantum transport in a quantum wire throughout the paper. In this section, we will first derive the eigenvalue equation, and then study the quantum transport and discuss how to manipulate the flying qubits in this system.

A. spin selective scattering

The total spin in \( z \) direction is always conserved during the time evolution of an arbitrary state. However, the applied magnetic field breaks the spin degeneracy. Therefore, dependent on the spin-orientation of both the Bloch electron and the magnetic impurity, the Hamiltonian in Eq. (1) can be rewritten as the direct sum of three parties: \( H_{ee}, H_{gg} \) and \( H_{eg} \). Here,

\[ H_{ee} = H_{qw}^{ee} + \Omega_l + \delta_e a_{0e}^\dagger a_{0e} \]  
\[ H_{gg} = H_{qw}^{eg} - \Omega_l + \delta_g a_{0g}^\dagger a_{0g} \]

is the Hamiltonian of the subspace with all spins down, where \( \delta_g = \omega_0 - \omega - \Delta - J_z \). The two-by-two matrix

\[ H_{eg} = \sum_{\sigma} H_{qw}^{e\sigma} + \frac{\omega_g}{2} a_{0g}^\dagger a_{0g} + \Omega_l \]

\[ -J_z a_{0e}^\dagger a_{0e} \]

\[ \bar{\omega}_e a_{0e}^\dagger a_{0e} - \Omega_l \]  

\[ \bar{\omega}_g a_{0g}^\dagger a_{0g} \]  

\[ -J_z a_{0g}^\dagger a_{0g} \]  

\[ \bar{\omega}_e a_{0e}^\dagger a_{0e} - \Omega_l \]  

\[ \bar{\omega}_g a_{0g}^\dagger a_{0g} \]  

\[ -J_z a_{0e}^\dagger a_{0e} \]

\[ \bar{\omega}_e a_{0e}^\dagger a_{0e} - \Omega_l \]  

\[ \bar{\omega}_g a_{0g}^\dagger a_{0g} \]  

\[ -J_z a_{0g}^\dagger a_{0g} \]

\[ \bar{\omega}_e a_{0e}^\dagger a_{0e} - \Omega_l \]  

\[ \bar{\omega}_g a_{0g}^\dagger a_{0g} \]  

\[ -J_z a_{0e}^\dagger a_{0e} \]

\[ \bar{\omega}_e a_{0e}^\dagger a_{0e} - \Omega_l \]  

\[ \bar{\omega}_g a_{0g}^\dagger a_{0g} \]  

\[ -J_z a_{0g}^\dagger a_{0g} \]

FIG. 2: (Color online). The reduced one-dimensional quantum wire: (a) in the subspace with all spins up or all spins down; (b) in the subspace with one spin up and one spin down.
describes the interaction in the subspace with one spin down and one spin up, where
\[
\bar{\omega}_g = J_z + \omega_0 - \omega - \Delta, \quad (6a)
\]
\[
\bar{\omega}_e = J_z + \omega_0 - \omega + \Delta. \quad (6b)
\]
In Eq. (5), the off-diagonal elements present the interaction of the electron spin with the x and y components of the impurity spins which leads to spin-flip. Obviously, the original single-band split into subbands with energy separations controlled by the local magnetic field applied to the impurity. In Fig. (2), we give a sketch of the equivalent quantum wire in different subspaces. In subspace with all spins up or down, it is a 1D chain with nearest interaction, however in other subspace corresponding to the mixture of the singlet and triplet states, it is two 1D chains which get crossed at the point \( j = 0 \).

### B. scattering in subspace with all spin down

Now we study the transport property of the Bloch electron in this system. Note that the configurations are similar in subspace \( \{ a^+_{jg} | 0g \} \) and \( \{ a^+_{je} | 0e \} \), as well as the subspace described by Hamiltonian \( H_{eg} \) when \( J_e = 0 \), so we take the scattering problem in the subspace with all spin down as an example. When the Bloch electron and the impurity are widely separated, we can separate the electronic and impurity-spin degrees of freedom. Therefore the eigenstate for this subspace is in the form
\[
|E_k\rangle = \sum_j u^+_k (j) a^+_{jg} |0g\rangle, \quad (7)
\]
where \( |g\rangle \) is the spin state of the magnetic impurity, \( u^+_k (j) \) is the probability density for finding the Bloch electron with spin down at the \( j \)th site. The wave function \( u^+_k (j) \) in this chain are obtained from the discrete Schrödinger equation
\[
(E_k - \epsilon_e - \delta_g \delta_{0j}) u^+_k (j) = \xi [u^+_k (j + 1) + u^+_k (j - 1)], \quad (8)
\]
where
\[
\epsilon_e = \omega - \Omega_J. \quad (9)
\]
If we regard the hopping between different sites as the kinetic term in Eq. (8), a delta-type potential is given rise to by the exchange interaction between spins and the on-site energy at site \( j = 0 \).

An incoming wave with energy \( E_k \), incident from the left, results in a reflected and transmitted wave. The wave functions in the asymptotic regions on the left and right are given by
\[
u^+_k (j) = \begin{cases} e^{ikj} + r_1 e^{-ikj} & j < 0 \\ t_1 e^{ikj} & j > 0 \end{cases}
\]

\[ (E - \epsilon_e + \delta_g) u^+_k (0) = \xi [u^+_k (1) + u^+_k (-1)] \quad (11b) \]

the transmission amplitude can be found
\[
t_1 = \frac{2\xi \sin k}{2\xi \sin k - \delta_g} \quad (12)
\]
and the backscattering amplitude also can be obtained by the relation \( t_1 = 1 + r_1 \), where the dispersion relation
\[
E_k = \epsilon_e + 2\xi \cos k \quad (13)
\]
is used.

As the singularities of the scattering matrix give rise to the resonant states, in Fig. 3 we plot the contour curve of the transmission coefficient \( T = |t_1|^2 \) in the complex plane. From Fig. 3(a), it can be seen that the transmission spectrum possesses two resonance poles in the complex-energy plane. Each pole is connected with a resonance level (or a quasibound state). According to the Breit-Wigner formalism, the real part of a pole presents the energy of a quasibound state, and the imaginary part is related to the lifetime of the resonance level \([20, 21]\). In Fig. 3(a), these two poles are separate on the real-energy axis, which means that each pole corresponds to a bound state. These states are produced by lack of periodicity in r-space. Bound states are the eigenfunctions in this system \([22]\). In this subspace with all spin down (or up), bound states has no contribution to the quantum transport, because scattering states survive only inside the band. However, in the subspace of antiparallel spins, shown in Fig. 2(b), things become different. We will discuss this situation in the following section. Obviously, the energy of these bound states is tunable with gate voltage at origin applied by a quantum point contact, and the exchange coupling along z direction as well as the magnetic field. Transmission zeroes also include in Fig. 3(a), which appear at the border of the band.
Transmission vanishes due to the vanishing group velocities at $k = 0, \pm \pi$. Figure 3(b) is the contour plot of transmittance as a function a complex momentum. It confirms the above results. Furthermore, transmission vanishes and two bound states appear outside the band at $\text{Re} k = 0, \pm \pi$. This result can be found analytically by substituting the complex momentum into the dispersion relation in Eq. (15b). Actually the energy and the formation of the bound states can be analytically gotten by the following assumption

$$u_k^j (j) = \begin{cases} C_1 e^{i(n\pi + \kappa)j} & j < 0 \\ C_1 e^{i(n\pi - \kappa)j} & j > 0 \end{cases},$$

where $n = 0, 1$ and $\kappa > 0$. From the continuous condition given in Eq. (11b), the energy of the bound state is obtained as

$$E_k = \omega + (-1)^n \sqrt{4\xi^2 + \delta_j^2},$$

and the normalized wavefunction reads

$$u_k^j (j) = \begin{cases} \sqrt{\tanh \kappa} e^{i(n\pi + \kappa)j} & j < 0 \\ \sqrt{\tanh \kappa} e^{i(n\pi - \kappa)j} & j > 0 \end{cases}.$$

Here, only bound states with even parity exist.

C. spin-dependent switch

We now investigate the scattering process of the propagating electron in a subspace with one spin up and one spin down. The stationary state in this excited space has the form

$$|E_k\rangle = \sum_j \left[ u_k^e(j) a_{jg}^\dagger |0e\rangle + u_k^g(j) a_{je}^\dagger |0g\rangle \right]$$

where $u_k^j (j)$ is the wave function of the Bloch electron with spin $i = e$ or $g$. Applying Eq. (17) into the stationary Schrödinger equation, one obtains

$$J_z \delta_{j0} u_k^e(j) = \xi \left[ u_k^e(j + 1) + u_k^e(j - 1) \right] - (E_k - \epsilon_e - \omega J_z) u_k^e(j)$$

$$J_z \delta_{j0} u_k^g(j) = \xi \left[ u_k^g(j + 1) + u_k^g(j - 1) \right] - (E_k - \epsilon_e - \omega J_z) u_k^g(j)$$

where $\omega_i (i = e, g)$ are given in Eq. (6), $\epsilon_e$ is listed in Eq. (9), the parameter

$$\epsilon_e = \omega + \Omega_I.$$

Equation (18) shows that there is a delta potential at $j = 0$ in each channel, which can be adjusted by on-site energy $\omega_0$, the Zeeman energy $\Delta$, and the exchange interaction strength $J_z$ along $z$ direction. Obviously, when $J_z = 0$, these two chains are independent. From the above discussion, we have already known that a discrete state is created by the impurity and this discrete energy is outside the band. In this section, new features arise due to the nonvanishing coupling $J_z$. It is well known that when a system is characterized by a coupling of a certain discrete energy and a continuum state, Fano resonance [20,23,24] appears because the discrete state offers one additional propagation path in the wave scattering which interact constructively or destructively.

A propagating electron in these channel will occupy an energy of the form

$$E_k = \epsilon_e + 2\xi \cos k_g = \epsilon_e + 2\xi \cos k_e,$$

where $k_e$ and $k_g$ are the electron wave vectors. The process that an incident wave impinges upon the structure under study, and transmitted and reflected wave emerge, is formulated by

$$u_k^e(j) = \begin{cases} e^{ik_ej} + r_k^e e^{-ik_ej} & j < 0 \\ t_{rl}^e e^{ik_ej} & j > 0 \end{cases},$$

$$u_k^g(j) = \begin{cases} D e^{-ik_ej} + r_k^g e^{ik_ej} & j < 0 \\ B e^{ik_ej} & j > 0 \end{cases}.$$

Applying Eq. (21) to the discrete Schrödinger equation (18) for the 0th and $\pm$1th sites, we immediately obtain the transmission amplitude

$$t_{rl}^e = \frac{2i\xi \sin k_g (2i\xi \sin k_e + \omega_e)}{(2i\xi \sin k_g + \omega_g) (2i\xi \sin k_e + \omega_e) - J_z^2},$$

within the channel, where the spins of the Bloch electron and the magnetic impurity are down and up respectively. For later convenience, we denote this channel as $|ge\rangle$ channel. The other channel in this subspace is denoted by $|eg\rangle$. The transmission amplitude from the $|ge\rangle$ channel to $|eg\rangle$ channel reads

$$B = \frac{2i\xi J_z \sin k_g}{(2i\xi \sin k_g + \omega_g) (2i\xi \sin k_e + \omega_e) - J_z^2}.$$

![Figure 4](image_url)
In the following discussion we set $n = 0$. The reflection amplitudes $r_\ell^g$ within the $|ge\rangle$ channel and $D$ from the $|ge\rangle$ channel to $|eg\rangle$ channel can be obtained by the relation $t_\ell^g = 1 + r_\ell^g$ and $D = B$ respectively. Here, we interest in a specific situation, namely, a regime where the Fano resonance might happen. Due to the removing the degeneracy of these two channels by the applied magnetic field, we can study the above regime by replacing $k_e$ with $n\pi + ik_e$. Then Eq. (22) becomes

$$t_\ell^g = \frac{2i\xi \sin k_g \left[\omega_e - (-1)^n 2\xi \sinh \kappa_e\right]}{(2\xi \sin k_g + \omega_g) \left[\omega_e - (-1)^n 2\xi \sinh \kappa_e\right]} - J_x^2$$

In the following discussion we set $n = 0$.

In Fig. 4 we plot the transmission coefficient $T_\ell^g = |t_\ell^g|^2$ (blue solid line), reflection coefficient $R_\ell^g = |r_\ell^g|^2$ (red dotted line), and the sum of $T_\ell^g + R_\ell^g$ (black dot-dashed line) as a function of the energy $E$. It can be found that: 1) there is a transmission zero (or a perfect reflection) inside the band; 2) the flow is preserved, which is implied by $T_\ell^g + R_\ell^g = 1$ in this regime. The vanishing transmission is caused by the coupling between the continuum and the bound state. The reason for the conserved flow is that except the bound state, other states are forbidden outside the band of $|eg\rangle$ channel, which is the energy regime we are interested, therefore energy is conserved in $|ge\rangle$ channel. The difference between Fig. 4a and others are that there is some overlap of the continuum between these two channels, which is why there are only two transmission zeroes in Fig. 4a and three transmission zeroes in other figures. From Eq. (18b), we know that there are no delta potentials at the 6th site in Fig. 4a, however there are a delta potential well in Fig. 4c) and a delta potential barrier in Fig. 4d). The well increase the transmission at the right side of the dip and decrease it at the left side of the dip, vice versa for barrier. Figure 5 shows the influence of the coupling strength $J_x$. It can be found that the width of the dip increases as $J_x$ increases. The above discussion implies that one can use the local magnetic field and gate voltage to control the transport of this flying qubit, a device with filtering and switching function is obtained.

IV. CONCLUSION

We have investigated the spin of an electron propagating along the host 1D structure, which interacts and exchanges information with a local magnetic impurity. By the tight-binding approach, the spin-dependent transmission/reflection coefficients is calculated within single-electron configurations. The evanescent states are found, due to the breaking periodicity of this 1D structure in r-space, and the energies of the evanescent states can be adjusted by the local electrode and the external magnetic field. In particular, we have shown that a total reflection occurs for a propagating electron spin incident with the energy of the evanescent state and anti-parallel to the spin orientation of the impurity. This total reflection is originated from the Fano resonance. Therefore, the Zeeman splitting and spin-flip interaction are necessary.

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[1] Y. L. Lim, A. Beige, and L. C. Kwek, Phys. Rev. Lett. 95, 030505 (2005).
[2] T. Durt, L. C. Kwek, and D. Kaszlikowski, Phys. Rev. A 77, 042318 (2008).
[3] J. Kim, O. Benson, H. Kan, and Y. Yamamoto, Nature (London) 397, 500 (1998).
[4] J.-T. Shen and S. Fan, Phys. Rev. Lett. 95, 213001 (2005).
[5] C.P. Sun, L. F. Wei, Yu-xi Liu, and F. Nori, Phys. Rev. A 73, 022318 (2006);
[6] P. Bermel, A. Rodriguez, S.G. Johnson, J.D. Joannopoulos, and M. Soljačić, Phys. Rev. A 74, 043818 (2006).
[7] D.E. Chang, A.S. Sørensen, E.A. Demler, and M.D. Lukin, Nat. Phys. 3, 807 (2007).
[8] Z. R. Gong, H. Ian, L. Zhou, and C.P. Sun, Phys. Rev. A 78, 053806 (2008).
[9] B. Dayan, A. S. Parkins, T. Aoki, E. P. Ostby, K. J. Vahala, and H. J. Kimble, Science 319, 1062 (2008).
[10] L. Zhou, Z.R. Gong, Yu-xi Liu, C.P. Sun, and F. Nori, Phys. Rev. Lett. 101, 100501 (2008).
[11] Lan Zhou, H. Dong, Yu-xi Liu, C. P. Sun and Franco Nori, Phys. Rev. A 78, 063827 (2008).
[12] S. Yang, Z. Song and C. P. Sun, Phys. Rev. A 73, 022317 (2006).
[13] D. Gunlycke, J. H. Jefferson, T. Rejec, A. Ramšak, D. G. Pettifor, and G. A. D. Briggs, J. Phys.: Condens. Matter
[14] J. H. Jefferson, A. Ramšak, and T. Rejec, Europhys. Lett. 74, 764 (2006).
[15] M. Habgood, J. H. Jefferson, A. Ramšak, D. G. Pettifor, and G. A. D. Briggs, Phys. Rev. B 77, 075337 (2008).
[16] D. Goldhaber-Gordon, H. Shtrikman, D. Mahalu, D. Abusch-Magder, U. Meirav, M.A. Kastner, Nature 391, 156 (1998).
[17] D. Goldhaber-Gordon, J. Gores, and M. A. Kastner, Phys. Rev. Lett. 81, 5225 (1998).
[18] N. Sundaram, S. A. Chalmers, P. F. Hopkins, and A. C. Gossard, Science 254, 1326 (1991).
[19] G. Binnig and H. Rohrer, Rev. Mod. Phys. 71, S324 (1999).
[20] A. M. Satanin and Y. S. Joe, Phys. Rev. B 71, 205417 (2005).
[21] Y.S. Joe, E.R. Hedin, and A.M. Satanin, Phys. Rev. B 76, 085419 (2007).
[22] N. Hatano, K. Sasada, H. Nakamura and T. Petrosky, e-print [arXiv:0705.1388v2].
[23] U. Fano, Phys. Rev. 124, 1866 (1961).
[24] J. Gores, D. Goldhaber-Gordon, S. Heemeyer, M. A. Kastner, H. Shtrikman, D. Mahalu, and U. Meirav, Phys. Rev. B 62, 2188 (2000).