Dynamical resurrection of the visibility in a Mach-Zehnder interferometer

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(Dated: January 15, 2013)

We study a single-electron pulse injected into the chiral edge-state of a quantum Hall device and subject to a capacitive Coulomb interaction. We find that the scattered multi-particle state remains unentangled and hence can be created itself by a suitable classical voltage-pulse \( V(t) \). The application of the inverse pulse \(-V(-t)\) corrects for the shake-up due to the interaction and rescues the original injected wave packet. We suggest an experiment with an asymmetric Mach-Zehnder interferometer where the application of such pulses manifests itself in an improved visibility.

PACS numbers: 73.23.-b, 03.65.Yz, 85.35.Ds, 73.43.Lp

On demand single-electron sources are an essential building block on the road to a mesoscopic solid state implementation of quantum computing. Single-particle wave packets can be generated with the help of suitable voltage pulses \[1\] and a first experimental realization of such a source has been recently achieved \[2\] in a quantum Hall setup. Contrary to their photonic counterparts, such single-electron states are prone to decoherence due to the interaction with the underlying Fermi-sea \[3\]. Here, we study the influence of a capacitive Coulomb interaction on a single-electron wave-packet injected into the chiral edge state of a quantum Hall device. Due to the interaction, the injected particle transfers energy to the Fermi sea, leading to the shake-up of electron-hole pairs. Analyzing the resulting scattered state, we find that it corresponds to a simple Slater determinant; the underlying product nature of the resulting multi-particle state allows one to undo the decoherence by applying a suitable local voltage-pulse.

The resurrection of decohered single-particle wave packets has numerous potential applications; here, we suggest to test this prediction in a Mach-Zehnder interferometer implemented in a quantum Hall setup. Electronic decoherence has become apparent in such devices \[4\] through the observation of a non-trivial decay of the visibility \[5, 6\] with increasing bias voltage and a satisfactory explanation could be obtained \[7\] via accounting for strong Coulomb interaction between edge states. The experiments \[4, 6\] have been performed with a finite bias voltage where electrons are stochastically injected into the system. We suggest to use an asymmetric setup operating in the \( \nu = 1 \) quantum Hall regime, where the decoherence is introduced in a controlled manner through a capacitive coupling of one arm to a metallic gate. Applying suitable voltage pulses to the scattered wave function behind the interaction region, the visibility of the interference pattern can be improved considerably though not perfectly, a consequence of our ignorance regarding the path which the electron has taken in traversing the device.

In the following, we study a one-dimensional ballistic conductor with chiral spinless electrons propagating to the right. The capacitive Coulomb interaction is described \[8\] by the Hamiltonian \( \hat{H}_{\text{int}} = \hbar \omega_c N^2/2 \), where \( N = \int dx g(x) :\hat{\Psi}(x)\hat{\Psi}(x)\) : is the effective number of excess electrons within the interaction region defined by the coordinate kernel \( g(x) \), \( \hbar \omega_c \) is the Coulomb energy, and \( :\hat{A}\hat{A}^\dagger:\equiv \hat{A} - \langle \hat{\Phi}_\nu | \hat{\Phi}_\nu \rangle \) denotes normal ordering with respect to the Fermi sea \( |\Phi_\nu\rangle \). Within our setup, we ignore any dissipative coupling to the environment \[3\] and study the effect of the interaction on the structure of the scattered many-body wave function.

Adopting a Luttinger Liquid description \[9\], we introduce the chiral bosonic field \( \hat{\theta}(x) = -\sum_{k>0}(\hat{b}_k e^{ikx} + \hat{b}_k^\dagger e^{-ikx})/\sqrt{2} \) obeying the commutation relation \( [\hat{\theta}(x), \hat{\theta}(x')] = i\pi \text{sgn}(x-x') \), where \( \hat{b}_k, \hat{b}_k^\dagger \) are bosonic creation and annihilation operators, \( [\hat{b}_k, \hat{b}_{k'}^\dagger] = \delta_{kk'} \). The electron field operator \( \hat{\Psi}(x) \) and the electronic density fluctuations \( \hat{\rho}(x) = :\hat{\Psi}(x)^\dagger\hat{\Psi}(x)\) : can be expressed via the field \( \hat{\theta}(x) \) as \( \hat{\Psi}(x) = (\sqrt{2}\pi \delta) \exp[-i\hat{\theta}(x)] \) and \( \hat{\rho}(x) = \partial_x \hat{\theta}(x)/2\pi \), where \( \bar{F} \) is the Klein factor acting as a Fermion-number ladder operator and \( \delta \) is an ultraviolet cutoff.

The interaction Hamiltonian is quadratic in the bosonic field \( \hat{\rho}(x) \), allowing for an exact solution of the equations of motion for \( \hat{\theta}(x, t) \), see Ref. \[2\]. Here, we choose a different approach and apply a Hubbard-Stratonovich transformation with an auxiliary real field \( z(t) \) in order to express the evolution operator \( \hat{S}(\rightarrow\infty, \infty) \) as an exponential linear in \( \hat{N}(t) \).

\[
\hat{S} = \hat{T}_+ \exp\left[ -i\omega_c \int dt \hat{N}^2(t) \right] \tag{1}
\]

\[
= \int D[z] \hat{T}_+ \exp\left[ i\omega_c \int dt (z^2(t) - z(t)\hat{N}(t)) \right],
\]

with \( \hat{T}_+ \) the usual (forward) time ordering operator. Below, we will be interested in a non-stationary situation and thus define the evolution operator \( \hat{S}_K \) as a time ordered exponent along the Keldysh contour, where \( z(t) \) assumes different values \( z_\pm(t) \) on the upper and lower branch of the contour \[10\]. Correlation functions are defined as \( \langle \hat{T}_K \{ \hat{S}_K \hat{A}(t_1^{\mu_1}) \hat{B}(t_2^{\mu_2}) \ldots \} \rangle \), where the Keldysh indices \( \mu_i \in \{ \pm \} \) specify the branch of the Keldysh con-
tour for the corresponding time instant $t$, and the average is
taken over the Fermi sea $|\Phi_F\rangle$.

We consider the scattering problem where an
electron wave packet $f(x)$ is injected at the left above
the Fermi sea, $|f\rangle = \int dx f(x-x_0)\tilde{\Psi}(t,x_0)|\Phi_F\rangle$, far
from the interaction region (we assume $t_0 \to -\infty$
and $x_0 = v_F t_0$, $v_F$ the Fermi velocity). At time $t = 0$, the
wave-packet reaches the interaction region where additional
electron-hole pairs can be excited. Finally, as $t \to +\infty$, the resulting many-particle scattering state $|\tilde{f}\rangle = \int dx f(x-x_0)\tilde{S}\tilde{\Psi}(t,x_0)|\Phi_F\rangle = \tilde{S}|f\rangle$ involves the
excess electron dressed by a cloud of electron-hole excitations propagating to the right.

Our focus is on the complexity of the scattered state $|\tilde{f}\rangle$. The simplest fermionic many-particle state is a
Slater determinant, i.e., an anti-symmetrized (A) product
state (obviously, our initial state $|f\rangle$ is of this type).
In fact, Slater determinants $A(\{|\phi_1\rangle \otimes |\phi_2\rangle \otimes \cdots \otimes |\phi_N\rangle\})$
correspond to non-entangled $N$-partite states which can be
created trivially by collecting fermionic particles in
quantum states $\{|\phi_i\rangle\}^N$ without use of any entanglement
resources, e.g., two-particle operations \[1\]. Furthermore, any two $N$-particle Slater determinants can be transformed
to one another by means of a suitable unitary operator $\tilde{U}$
generated by a single-particle Hamiltonian.

The scattered state is obtained from a unitary evolu-
tion which involves two-particle interactions and thus $|\tilde{f}\rangle$ in general is not expected to be a simple Slater
determinant, even if the incoming state is one. Below, we
analyze the complexity of $|\tilde{f}\rangle$ by testing its overlap with
any Slater determinant $|f_U\rangle \equiv \tilde{U}|f\rangle$ which we evolve from
the initial state $|f\rangle$ via a unitary single-particle operator
$\tilde{U}$ \[12\]. Maximizing the overlap $\langle f_U|\tilde{f}\rangle$ as a function of
the evolution operator $\tilde{U}$, we find the optimal operator
$\tilde{U}_{\text{opt}}$ providing the Slater determinant $|f_{U_{\text{opt}}}\rangle$ closest to
$|\tilde{f}\rangle$. For the chiral electrons discussed here, we find that
the scattered state is again a Slater determinant.

We restrict ourselves to physically relevant single-
particle operators $\tilde{U}$ that can be generated through ap-
lication of a local voltage-pulse $V(t)$ at the position $x$
of the wire. The operator

$$\tilde{U}_x = \tilde{T}_+ \exp \left[ i v_F \int dt \chi(t) \tilde{\rho}(x,t) \right] \quad (2)$$

then adds the phase $\chi(t) = (e/\hbar) \int dt' V(t')$ to each elec-
tron passing the position $x$. Furthermore, we change perspec-
tive and rotate the scattered state $|\tilde{f}\rangle$ by $\tilde{U}_x$. In
maximizing the overlap $\langle f|\tilde{U}_x|\tilde{f}\rangle$, one then has to find a
pulse $V(t)$ such as to bring the resulting state as close as
possible to the original incoming state $|f\rangle$. The overlap
$\langle f|\tilde{U}_x|\tilde{f}\rangle$ can be expressed as

$$\langle f|\tilde{U}_x|\tilde{f}\rangle = \int dxdx' f_2(x)f_1(x') \times \langle \tilde{T}_K\{\hat{S}_K \tilde{U}_x \hat{\Psi}(x,t_{+2}^T)\hat{\chi}(x',t_{+2}^T)\} \rangle. \quad (3)$$

The wave packets $f_{1,2}(x) = f(x-x_{1,2})$ are localized
around the retarded coordinates $x_{1,2} = v_F t_{1,2}$ far left
(right) to the interaction region (thus rendering the over-
lap independent on the time coordinates) and the uni-
tary operator $\tilde{U}_{xK}$ has a nonvanishing phase $\chi(t)$ only
on the upper branch of the Keldysh contour (hence
$\chi(t) = (t_{+T})$). The average in Eq. (3) involves a product
of exponentials linear in the bosonic fields; in addition,
we choose an incoming wave packet of Lorentzian form $f_1(x)
= \sqrt{\xi/\pi} (x+i\xi)^{-1}$ with width $\xi$. A straightforward
calculation provides the (zero-temperature) result

$$\langle f|\tilde{U}_x|\tilde{f}\rangle = \exp \left[ -\int_0^\infty \frac{d\omega}{2\pi} \left( \frac{\omega^2}{v_F^2} + 2\text{Re}(\chi(\omega))\bigg[ e^{-\omega \tau_1} + \frac{i\omega v_F}{2} \right) \left( e^{-\omega \tau_1} - \frac{\omega v_F}{2} \bigg) \right] \right], \quad (4)$$

with the time-width $\tau_1 = \xi/v_F$ and the transformed in-
teraction kernel $g(\omega) = \int dt g(v_F t) e^{-i\omega t}$ (we assume $g(x)$
to be centered around $x = 0$). $G_{++}(\omega)$ is the Fourier
transform of the bosonic Green’s function $G_{++}(\omega, x, x) =
\langle \tilde{T}_+ \{ \tilde{\rho}(x, \tau) \tilde{\rho}(x, 0) \} \rangle$,

$$G_{++}(\omega, x, x') = \frac{1}{2\pi v_F^2} \int \frac{d\omega'}{2\pi} e^{-i\omega'(x-x')/v_F} \frac{g(\omega')}{\omega' - \omega - i\delta \text{sgn}(\omega)}, \quad (5)$$

and $\Pi_{++}(\omega)$ derives from $\Pi_{++}(\tau) = \langle \tilde{T}_+ \{ \tilde{N}(\tau) \tilde{N}(0) \} \rangle$, $\Pi_{++}(\omega) = \int dx dx' g(x) g(x') G_{++}(\omega, x, x')$. Maximizing
the overlap $\langle f_U|\tilde{U}_x|\tilde{f}\rangle$ with respect to the phase $\chi(t)$ one
finds that the optimal voltage-pulse must generate the
phase

$$\chi_L(\omega > 0) = -\frac{\omega v_F}{2} \frac{|g(\omega)|^2 e^{-\omega \tau_1}}{1 - i\omega v_F \Pi_{++}(\omega)/2}, \quad (6)$$

and $\chi_L(\omega < 0) = \chi_L^*(\omega)$.

Calculating the overlap $\langle f_U|\tilde{U}_x|\tilde{f}\rangle$ for this pulse one
arrives at the result that the scattered state is a Slater
determinant state, $\langle f_U|\tilde{U}_x|\tilde{f}\rangle = 1$. Hence, the (prop-
erately retarded) voltage-pulse generating the phase $\chi_L(t)$
rotates the scattered state $|\tilde{f}\rangle$ exactly back to the original
state, $\tilde{U}_{x\tilde{f}}|\tilde{f}\rangle \propto |\tilde{f}\rangle$, up to a phase factor.

The incoming state $|f_U\rangle$ with a Lorentzian wave packet
above the Fermi sea can itself be obtained \[1\] by applying a voltage pulse of Lorentzian shape \(V_i(t) = (2\hbar/e\tau)\left(1 + (t/\tau)^2\right)^{1/2}\) with the associated phase \(\tilde{\phi}(t) = 2\arctan(t/\tau)\), i.e., \(|f_i\rangle = U_{\phi(t)}(\hat{F}^\dagger)|\Phi_f\rangle\). The operator \(U_{\phi(t)}\) alone cannot add charge to the system and it is the Klein factor that adds an additional electron at the Fermi level.

Let us then apply an arbitrary voltage pulse \(V(t)\) generating an electron-hole state \(|f_0\rangle = U_{\phi(t)}(\hat{F}^\dagger)^n|\Phi_f\rangle\) with \(n\) excess electrons above the Fermi sea. Although not covering all possible incoming states, this procedure generates an important family of states which can be classically prepared by an experimenter. The action of the Coulomb interaction on this incoming state can then be expressed through a change in the phase \(\phi\), i.e., the scattered state \(\tilde{S}|f_0\rangle\) is equivalent to the state \(\hat{U}_G(\hat{F}^\dagger)^n|\Phi_f\rangle\) with the phase

\[
\hat{\phi}(\omega > 0) = \phi(\omega) \frac{1 + i\omega_c \Pi_{++}(\omega)/2}{1 - i\omega_c \Pi_{++}(\omega)/2}.
\] (7)

Note that it is the rotation by \(\hat{U}_{\phi-\phi} = \hat{U}_\chi\) that plays the role of the optimal rotation Eq. \(2\).

On a technical level, the reason for the separability of the scattered state is found in the chiral nature of the scattering problem. In the bosonic language, the incoming state can be presented as a coherent state \(|f_i\rangle = \int f_{x>0} e^{i u_x b_x - u_x^* b_x^\dagger} \hat{F}^\dagger|\text{vac}\rangle\), with \(u_k = v_F \sqrt{\phi(0)} (k v_F) e^{i k x}/2\pi\) and \(|\text{vac}\rangle\) the bosonic vacuum. Since the bosonic Hamiltonian is quadratic and in the absence of back reflection, the bosons only acquire a phase factor during the propagation through the interaction region, \(b_k \rightarrow e^{i k x} b_k\). Thus the scattered state is also a coherent state with \(u_k \rightarrow u_k e^{i k x}\) and therefore exhibits the same complexity as the incoming state.

We now make use of the results Eqs. \(6\) and \(7\) in a specific application, the improvement of the visibility in a Mach-Zehnder interferometer (MZI) fed through single-electron pulses. The basic idea is to undo the distortion of the wave packets by the Coulomb-interaction via suitable voltage-pulses applied after the scattering region. Consider an asymmetric MZI where the electrons passing through one arm (the bottom arm d, see Fig. \(\ref{fig:1}\)) are capacitively coupled to a metallic gate. We inject a single-electron pulse generating the phase \(\chi(t)\) is applied right after the interaction region,

\[
I_4(x_i) = \frac{e v_F}{\pi} \int \left|t_{ud}(x_i)^2 + \left|t_u(x_i)^2\right| f(x_i)^2 + t_u^* t_d f(x_i) G_\chi(x_i)\right| dx_i + \text{c.c.}
\]

where \(x_i = x - v_F t_i\), \(t_{ud}(x)\) are amplitudes to go from lead 2 to lead 4 through the upper (lower) arm and

\[
G_\chi(x_i) = \int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle,
\]

\[
I_4(x_i) = \int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle,
\]

where \(x_i = x - v_F t_i\), \(t_{ud}(x)\) are amplitudes to go from lead 2 to lead 4 through the upper (lower) arm and

\[
G_\chi(x_i) = \int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle,
\]

\[
I_4(x_i) = \frac{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle}{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle},
\]

\[
I_4(x_i) = \frac{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle}{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle},
\]

\[
I_4(x_i) = \frac{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle}{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle},
\]

\[
I_4(x_i) = \frac{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle}{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle},
\]

\[
I_4(x_i) = \frac{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle}{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle},
\]

\[
I_4(x_i) = \frac{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle}{\int dx' f^*(x' - x_0) \langle \tilde{T}_K \{ \tilde{S}_K \hat{\Psi}_d(x', t_0) \hat{\Psi}_d^\dagger(x, t)}\rangle},
\]
with \( \hat{\Psi}_{f_k}^+ = \int dx f_k(x) \hat{\Psi}^+(x) \). For a compensation \( \chi_{Mz} \), the visibility is given by the overlap of the state \( \hat{\Psi}_{f_k}^+ \hat{\Upsilon}_{Mz}(|\Phi_f\rangle \) (first term in \(|A\rangle \); the particle enters the lead 4 from the upper arm) and the contribution \( \hat{\Upsilon}_{Mz} \hat{\Upsilon}_{-Mz}(|\Phi_f\rangle \) (second term in \(|A\rangle \) with the particle coming from the lower arm). The different shapes of these wave functions provide `which-path' information which implies a reduction in the visibility. If the electron trajectory were known, one could fully compensate the effect of the interaction by applying \( \chi_{Mz} = \chi_L/e \) every time when the electron chooses the lower arm. Since this quantum information is not available, the pulse has to be applied blindly and either compensates the effect of interaction (if the particle indeed passed through the lower arm) or creates additional electron-hole pairs (if the particle passed through the upper arm). It turns out, that a half-pulse \( \chi_{Mz} = \chi_L/2 \) partly erases the information on the trajectory taken by the electron: after recombination in the second beam splitter, the many-particle states in the outgoing lead corresponding to the propagation through the upper or the lower arm are optimally adjusted to each other such as to increase the visibility. Or, in other words, the half-pulse \( \chi_L/2 \) on purpose erases part of the `which-path' information in the outgoing lead.

In reality the form of the coordinate kernel \( g(x) \) is not known. However, one can measure the time dependence of the function \( \Phi(t) \) which implies a reduction in the visibility. If the electron trajectory were known, one could fully compensate the effect of interaction by applying \( \chi_{Mz} = \chi_L/e \) every time when the electron chooses the lower arm. Since this quantum information is not available, the pulse has to be applied blindly and either compensates the effect of interaction (if the particle indeed passed through the lower arm) or creates additional electron-hole pairs (if the particle passed through the upper arm). It turns out, that a half-pulse \( \chi_{Mz} = \chi_L/2 \) partly erases the information on the trajectory taken by the electron: after recombination in the second beam splitter, the many-particle states in the outgoing lead corresponding to the propagation through the upper or the lower arm are optimally adjusted to each other such as to increase the visibility. Or, in other words, the half-pulse \( \chi_L/2 \) on purpose erases part of the `which-path' information in the outgoing lead.

We close with the explicit calculation of the visibility and its resurrection for a model kernel \( g(x) = \exp(-|x|/a) \) and an incoming Lorentzian. The optimal voltage-pulse \( V_{\text{c}}(t) = e_{\text{opt}} \chi_{Mz}(t)/\hbar \) derives from the phase

\[
\chi_{Mz}(t) = -4\hbar^2 \int_0^{\tau} d\nu \frac{e^{i\nu/\tau} - |\nu|/n_f}{(\nu + i)^2[(\nu - i)^2 - n_f^2].}
\]

With \( n = \alpha \rho \delta \varepsilon \) the number of electrons in the interaction region \( A \) within the energy window \( \delta \varepsilon \) (\( \rho = 1/\hbar v_F \) denotes the density of states), we can define the two parameters \( n_f = a/\xi \) and \( n_c = \omega_c \tau/4\pi \) with the corresponding energy scales \( \delta \varepsilon_f = hv_F/\xi \) associated with the excess energy carried by the wave packet and \( \delta \varepsilon_c = \hbar \omega_c/\xi \) associated with the typical energy of electron-hole excitations (electron-hole pairs with higher energies are suppressed). The (un)corrected visibilities \( (V_i)_{\text{opt}} \) are given by

\[
V_i = \exp \left[ 2i\hbar n_c \int_0^{\tau} d\nu \frac{\text{sgn}(\nu) e^{-2|\nu|/n_f}}{(\nu + i)^2[(\nu - i)^2 - n_f^2]} \right], \quad (\text{un})\text{corrected visibility.}
\]

and are shown in Fig. III as a function of interaction strength \( n_c \) for high- and low-energy incoming wave packets. The additional voltage-pulse indeed improves the visibility in both cases and the correction is more efficient at large \( n_c \), where the visibility saturates for fixed \( n_f \). The saturation at large Coulomb energy occurs due to energy conservation: an incoming electron cannot excite electron-hole pairs with energy higher than \( \delta \varepsilon_f \) no matter how strong the Coulomb energy is. The enhancement in visibility is more pronounced at large \( n_c \) and can reach of order 100\% of the original `bare' value. Summarizing, we have established the product nature of single-particle wave packets decohered through capacitive Coulomb interaction and have demonstrated the possibility for their resurrection through appropriate voltage pulses. In a real device, the interaction may add dissipation to the system [3], which our scheme cannot cure. Nevertheless, those parts of decoherence/dissipation which are due to particle shake-up (i.e., a deformation of the wave-packet) can always be compensated by application of proper voltage pulses.

We thank Gordey Lesovik for discussions and acknowledge the financial support from the Swiss National Foundation through the Pauli Center at ETH Zurich.

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