Localization of Quantum States at the Cyclotron Resonance.

V. Ya. Demikhovskii, D. I. Kamenev

Nizhny Novgorod State University, Gagarin ave. 23, Nizhny Novgorod, 603600, Russia

E-mail: demi@phys.unn.runnet.ru

Tunneling of a quantum particle to the classically inaccessible region in the intrinsically degenerated system is investigated here by means of quasienergy eigenstates. The exact resonance $\delta \omega = \omega - \omega_c = 0$ and near-resonance $\delta \omega \neq 0$ cases are explored numerically. It is shown that in both cases all quantum states are localized. Correspondence of the quantum dynamical barriers to the classical invariant curves is demonstrated. The phenomena considered in this article can be observed when an ultrasound wave propagates perpendicularly to a magnetic field and interacts with a 2D electron gas in a semiconductor heterostructure.

In this paper we treat the role of quantum dynamical barriers in localization of quantum states at the cyclotron resonance. Dynamical barriers in quantum systems are counterparts of classical invariant tori (or KAM tori) which block diffusion of a particle along the phase space. Similarly, in quantum case dynamical barriers inhibit the diffusion of a wave packet in a Hilbert space that leads to localization of quantum states. This phenomenon was previously analysed for the kicked rotator model [1, 2] and the nonlinear oscillator in a monochromatic radiation field [3]. However quantum dynamics was explored mainly for the systems, whose classical counterparts are accidental degenerate and the conditions of the KAM theorem are satisfied. The dynamics of the intrinsic degenerate system (kicked simple harmonic oscillator) was explored by Berman, Rubaev, Zaslavsky [5] under the condition of strong and weak chaos of the classical limit. In this work by using quasienergy eigenstates we treat tunneling through the dynamical barriers and localization of quantum states in an intrinsically degenerated system where the KAM theorem is invalid.

Our system is a charged particle subject to both a uniform magnetic field and a field
of a longitudinal monochromatic wave, propagating perpendicularly to the magnetic field direction. The phase space of the classical system at the condition of the cyclotron resonance \( \omega = \omega_c \), where \( \omega \) and \( \omega_c \) are the wave and cyclotron frequencies, consists of the infinit number of cells comprising closed classical orbits and separated by a separatrix lattice \([4]\). In the classical case in the resonance approximation a particle has no possibility to penetrate through a separatrix from one cell to another. In this paper we will show that in the quantum case the dynamical barriers become transparent for the wave packet even in the resonance approximation.

The theory of nonlinear resonance in the intrinsically degenerated system was developed by G.A.Luna-Acosta and the authors in Ref. \([6]\). It was shown that the Hilbert space of the discussed system breaks up into quantum resonance cells. The boundaries of these cells in the quasiclassical limit correspond to the the separatrices in the classical phase space. In this paper we numerically investigate the role of these ”quantum separatrices” in localization and tunneling phenomena at the cyclotron resonance.

The Hamiltonian of the considered system has the form

\[
\hat{H} = \left( \frac{\hat{p} - \frac{e}{c} \mathbf{A}}{2m} \right)^2 + V_0 \cos(kx - \omega t) = \hat{H}_0 + V(x, t),
\]

(1)

where \( e \) and \( m \) are the charge and mass of a particle, \( c \) is the light velocity, \( \hat{p} \) is the momentum, \( k \) and \( V_0 \) are the wave vector and amplitude of the perturbation. It is convenient to choose the gauge of \( \mathbf{A} \) in the form \( \mathbf{A} = (0, Hx, 0) \) in order to produce the magnetic field \( \mathbf{H} \) along the \( z \)-direction. In this gauge the momentum \( p_y \) is an integral of motion, hence we have to determine the dependence of the wave function only on two variables, \( x \) and \( t \).

Since the perturbation is periodic in time, the Floquet theory can be used for description of time evolution of the system in terms of quasienergy (QE) spectra and QE eigenfunctions \( \psi_q(x, t) \). The QE eigenfunctions are eigenstates of the evolution operator for one period of oscillation of the external field \( T = 2\pi/\omega \). They can be defined by (see for example Ref. \([7]\))

\[
\psi_q(x, t) = \exp\left(-\frac{iE_q t}{\hbar}\right) \sum_n \sum_s A_{n,s}^q \psi_n(x) \exp(-is\omega t),
\]

(2)
where $E_q$ is the QE eigenvalues, $\psi_n(x)$ is the nth eigenfunction of a simple harmonic oscillator, $\hbar$ is the Plank’s constant. After substituting Eq. (2) into the time-dependent Schrödinger equation, we obtain the system of the uniform algebraic equations:

$$
(E_q - \hbar\omega_c n + \hbar\omega s)A_{n,s}^q = \sum_m V_{n,m}(A_{m,s+1}^q + A_{m,s-1}^q),
$$

(3)

where the quasienergy $E_q$ is counted from the ground state $\hbar\omega_c/2$, the matrix elements $V_{n,m}$ was defined in Ref. [6].

First let us consider the exact resonance case where $\delta\omega = \omega_c - \omega = 0$. If $V_0 = 0$ each definite value $n - s$ in Eq. (3) corresponds to an infinitely degenerated QE level and these levels are separated by the distance $\hbar\omega$. When $V_0 \neq 0$ the degeneracy is broken and each initially degenerated level splits into QE spectrum. All the spectra are identical, and the distance between the extreme levels in one spectrum has the order of $V_0$. If $V_0 \ll \hbar\omega$ we can, at first order, neglect the interaction between the individual spectra and treat one of them.

QE eigenvalues $E_q$ and eigenfunctions $A_{n,s}^q$ were calculated in Ref. [6] for resonant ($\omega = \omega_c$) and near resonant ($\delta\omega \neq 0$) case by the perturbation theory assuming $V_0/\hbar\omega << 1$ and $\hbar\delta\omega n \ll \hbar\omega$ for all studied values of $n$. In the first order approximation Eq. (3) yields

$$
(E_q - \hbar\delta\omega n)A_n^q = V_{n,n+1}A_{n+1}^q + V_{n,n-1}A_{n-1}^q,
$$

(4)

where $A_n^q \equiv A_{n,n}^q$. The matrix elements $V_{n,n\pm1}$ are symmetric, real and for $n \gg 1$ can be expressed via the Bessel functions of the first order

$$
V_{n,n+1} = \frac{V_0}{2} \sqrt{\frac{n}{n+1}} e^{-\frac{\hbar}{2}} J_1(\sqrt{2n\hbar}),
$$

(5)

where $\hbar = (ka)^2$, $a = (\hbar \kappa^2)^{1/2}$ is the magnetic length. When $\delta\omega = 0$, the amplitude of the wave $V_0$, which is measured in units of $\hbar\omega$, does not influence the eigenvectors $A_n^q$ but changes only the energy scale, hence the effective Plank’s constant $\hbar$ is the only parameter in this problem. The oscillating matrix elements $V_{n,n+1}$ define the quantum resonant cells — relatively independent dynamical regions in the Hilbert space [6]. Each quantum cell corresponds to two classical cells (in action-angle variables).

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When $\delta \omega \neq 0$ and $V_0 = 0$ according to Eq. (4) the levels are separated from each other by the distance $\hbar \delta \omega$. If $V_0 \neq 0$ interaction between the levels appears, however, due to Eq. (4) this interaction is essential only within the region $n\hbar \delta \omega \sim V_0$, whereas the levels satisfied the condition $n\hbar \delta \omega \gg V_0$ practically are not affected by the perturbation.

Let us discuss the main features of the QE spectrum which substantially define dynamics of the system. In the exact resonance case the spectrum $E_q$ is almost equidistant near the top and the bottom at the distance $\hbar \tilde{\omega}$ [6]. The frequency $\tilde{\omega}$ in the quasiclassical limit is identical to the frequency of small oscillations near the center of the resonance in the phase space. In order to derive $\tilde{\omega}$ for the near resonance case, let us expand $f(n) \equiv V_{n,n+1}$ near some point $n_0$, defined below, into Taylor series and introduce the differential operator by

$$A_n \equiv \frac{d^2}{dn^2}A_n + 2A_n.$$ (6)

Then the finite-difference Shr"{o}dinger equation (4) takes the form

$$[E_q - (n - n_0)(\hbar \delta \omega + 2f'(n_0)) - 2f(n_0)]A_n = f(n_0)\frac{d^2}{dn^2}A_n + f''(n_0)(n - n_0)^2 A_n,$$ (7)

where $E_q = E_q - \hbar \delta \omega n_0$. If the condition

$$\hbar \delta \omega = -2f'(n_0)$$ (8)

is satisfied equation (7) takes the simple harmonic oscillator form with the center of oscillations $n_0$ defined by Eq. (8). According to Eq. (8), when $\delta \omega \neq 0$ there is only a finite number of the resonances. If $V_0$ is small, i.e. $V_0 < \hbar \delta \omega$, the unperturbed levels $n\hbar \delta \omega$ are not coupled by the perturbation (see Eqs. (4) or (8)) and the resonance disappears. The same is true for the classical case. In the limit $\hbar \to 0$, $n \to \infty$, and finite $n\hbar = I$, the positions of the resonances and the frequency, defined by Eqs. (8) and (7), correspond to the fixed points and frequency of small oscillations near the center of the resonance in the phase space [8].

An arbitrary solution of the time-dependent Schrödinger equation with the Hamiltonian (1) can be expanded in the simple harmonic oscillator basis

$$\psi(x, t) = \sum_n C_n(t)\psi_n(x)exp(-iE_n t/\hbar),$$ (9)
were $E_n$ is the energy of nth Landau level. Coefficients $C_n(t)$ determine the time development of the probability distribution defined by $|C_n(t)|^2$. Instead of integration of the time-dependent Schrödinger equation we express the coefficients $C_n(t)$ in terms of the Green function $G_{n,n'}$

$$C_n(t) = \sum_{n'} G_{n,n'}(t, t_0) C_{n'}(t_0),$$

which can be determined by using the complete set of the QE eigenstates

$$G_{n,n'}(t, t_0) = \sum_q A^q_n A^q_{n'} e^{-iE_q(t-t_0)/\hbar}.$$  

Thus, we have the following procedure for computing the probability distribution at time $t$ which is used below. First, by numerical solving the finite-difference Schrödinger equation (4) we compute the QE eigenfunctions $A^q_n$ and the QE eigenvalues $E_q$, then using the QE eigenstates we compute the Green function (11), and finally we use the Green function for calculating coefficients $C_n(t)$. Since the dependence of the coefficients $C_n(t)$ on time $t$ in Eqs. (10), (11) is explicit, they can be easily computed at arbitrary time $t$.

We have performed the numerical experiments in order to explore how the dynamical barriers in the intrinsically degenerated system (web tori [9]) influence the evolution of quantum states with various initial conditions at the cyclotron resonance. The system consists of the weakly interacting quantum resonant cells [6] defined by the matrix elements $V_{n,n+1}$. Our purpose is to treat electron tunneling between the cells.

As mentioned above, the QE states completely define the dynamics of the system. If the initial state is any QE eigenstate, i.e. $C_{n'}(t_0) = A^q_{n'}$, then in accordance with Eqs. (10) and (11) the probability distribution does not depend on time $t$. In particular, if the initial wave function $A^q_{n'}$ is localized in one cell then it will be localized in the initial cell for any time $t$. The localized QE eigenfunctions correspond to the eigenvalues near the top and bottom of the QE spectrum. By employing a coherent state representation (Husimi function) we can show that the localization of a wave packet in the Hilbert space entails the localization of the Husimi function in the phase space. The contour plots of two Husimi functions
corresponding to the extreme eigenvalues from the top and bottom of the QE spectrum are plotted in Fig. 1 (a) (the upper part of the figure corresponds to the top eigenvalue and the lower part corresponds to the bottom one). The symmetry of these Husimi functions can be explained from Eq. (4) which is invariant under the transformation

$$E_q \rightarrow -E_q, \quad A_n^q \rightarrow (-1)^n A_n^q,$$

(12)
corresponding to the transformation $x \rightarrow -x$ (or $\omega t \rightarrow \omega t + \pi$) in Eq. (2). The contour plots of the Husimi functions correspond to the classical orbits in two cells of the phase space which are shown in Fig. 1 (b). Another kind of initial state is an excited Landau level $C_n(0) = \delta_{n_n0}$. In this situation, the initial wave packet contains all the QE eigenstates and one can expect that such a state will be extending most intensively, tunneling to other resonant cells.

In our calculations the system involved 432 Landau levels which formed 7 resonance cells in the Hilbert space, shown in the upper part of Fig. 2. The initial state in the form $C_n(0) = \delta_{n_n0}$ was placed in the center of the first cell. After the time $\tilde{T} \sim 2\pi/\tilde{\omega}$ the wave packet spreaded over the initial cell. In parallel with the fast dynamics within the initial cell a slow process of propagating the probability distribution to the subsequent cells took place. The characteristic time of the probability distribution spreading over all the considered cells has the order of $2\pi/\omega_{min}$, where $\omega_{min}$ is the minimal distance between the QE eigenvalues in the QE spectrum calculated for seven cells. Two snapshots of evolution of the probability distribution $|C_n|^2$ (in a logarithmic scale) as a function of the Landau number $n$ are shown in Figs. 2 (a), (b). It is seen that the probability distribution in average decreases exponentially, but sharp decay occurs only at the boundaries of the cells that verifies our assumption of the role of these boundaries as dynamical barriers to the probability flow.

The probability distribution successively penetrates through the dynamical barriers from one cell to the subsequent ones, travelling along the Hilbert space. After filling some cell the wave packet reflects from the barrier and interferes (in the sixth cell in Fig. 2 at the moment (a) $t = 4 \times 10^5$ and (b) $t = 7 \times 10^5$, time is measured in units of $T = 2\pi/\omega$).
As a result, oscillations of the probability distribution appear. These oscillations can be explained also from another point of view. Let us consider tunneling of some initial wave packet, concentrated on a level with Landau number $n$ in an initial cell, to another level with number $n'$, situated in another cell. The main contribution to the Green function (11) is given by the QE eigenfunctions with those numbers $q$, for which both $A^n_q$ and $A^{n'}_q$ are sufficiently large or, in other words, any of these QE eigenfunctions must occupy two (or more) resonant cells. Such eigenvectors correspond to the eigenvalues near the center of the QE spectrum, i.e. the point $E_q = 0$ [6]. As follows from our calculations, the most delocalized QE eigenfunctions at the same time turn out to be most oscillating (by module) functions of $n$. The probability distribution in Fig. 2 also oscillates with the minimal period.

It is interesting that the probability distribution in average is greater near the boundaries of the cells (by 1.5 - 3 orders in Fig. 2). Since these boundaries correspond to the separatrices [6] (web tori), this quantum phenomenon is akin to the diffusion of a classical particle along the separatrice lattice within exponentially small stochastic regions in the phase space when nonresonant terms in the classical Hamiltonian are taken into consideration [9].

The probability distribution in Fig. 2 is approximately the same within a cell, hence it is reasonable to consider average probabilities, defined by $P_i(t) = \sum n_i |C_{n_i}(t)|^2$, where $n_i$ takes the values within the $i$th resonant cell. The evolution of $P_1(t)$ and $P_2(t)$ (in the first and second cells) is presented in Fig. 3 (the time axis is plotted in a logarithmic scale).

For the calculations we used the system with the same parameters as in Fig. 2 involving seven resonance cells. We suppose that the first two cells do not "feel" the limited boundary conditions because the probability distribution at the boundary of the system (the seventh cell) is exponentially small and including of more number of cells does not affect $P_1(t)$ and $P_2(t)$. The probability distribution in Fig. 3 does not leak out to the subsequent cells in the limit $t \to \infty$, hence we can conclude that the quantum states are localized. The localization of the quantum states suggests a conjecture (Ref. [6]) about discreteness of the QE spectrum for the unbounded system when $n \to \infty$. Time-averaged probabilities $\overline{P_i}(t)$ versus cell number for two values of the effective Plank's constant $\hbar$ are presented.
in Fig. 4. In the logarithmic scale the exponential behaviour of these functions is fairly evident. As one can see in Fig. 4 the relation \( \frac{P_2(t)}{P_1(t)} \) has the order of \( 10^{-2} \). It should be stressed that the sharp decay of the probability distribution occurs not everywhere, but only at the boundaries of the quantum resonance cells. Thus, we deal with the new kind of localization, namely, localization over the quantum resonant cells. The localization length in the resonance approximation does not depend on the amplitude of the wave \( V_0 \). It is defined by the only parameter \( h = (ka)^2 \) which determines the behaviour of the Bessel function in the matrix elements in Eq. (5). Our numerical experiments have shown that slope of the curve on Fig. 4, defining the localization length, in general increases with increasing \( h \). However, the decay of the probability distribution at different boundaries between the cells changes randomly. The nature of this randomness is discussed below.

Next effect, which we consider here, follows from discreteness of the quantum number \( n \) which labels Landau states in the Hilbert space. If we choose the parameter \( h \) so that

\[
V_{n_0,n_0+1} \sim J_1(\sqrt{2n_0h}) = 0
\]  

(13)

then in accordance with Eq. (4) transitions through the level \( n_0 \) are blocked and the Hilbert space is divided into two disjoint parts. If we will change the parameter \( h \), the condition (13) will be successively satisfied for the \( n_0 + 1 \)-th, \( n_0 + 2 \)-th, \ldots levels. Thus the tunneling probability through the dynamical barrier goes periodically to zero. Let us define the penetration coefficient by

\[
P = \sum_{n>n_0} |C_n(T)|^2,
\]

where \( n_0 \) is located at the boundary between the second and the third cells, and place the initial state in the form \( C_n(0) = \delta_{n,n'} \) at the center of the second cell. The oscillations of \( P \) with changing the parameter \( 1/h \sim H \) are shown in Fig. 5, where the above-mentioned zeroth values of the penetration coefficient have been replaced by small numbers. The oscillations have the period \( \Delta H = 2hck^2/eb^2 \), where \( b \) is a root (the second one in our case) of the Bessel function (5). The amplitude of the oscillations decreases with decreasing \( h \), that does not contradict to the quasiclassical formula \( P \sim \exp(-\alpha/h) \). The definition of more explicit form of the dependence of the tunneling probability \( P \) on \( h \) is complicated by the wild oscillations of the dynamical barrier.
permeabilities as a function of $h$.

The dynamical barriers become impermeable for definite values of $h$ only in the resonance approximation when the condition $\tilde{\omega} \ll \omega$ is satisfied and we take into account only transitions between the nearest levels.

Next we present the results of calculations for the near resonance case $\delta \omega \neq 0$ which are illustrated in Fig. 6 where the probability distribution spreading after the time $t = 10^6$ is shown. There is a significant difference between the near resonance and exact resonance cases. Though the matrix elements (5) are the same as in the exact resonance case, the additional term $n\hbar \delta \omega$ in (4) deforms the resonance cells with small $n$’s and destroys the cells with large $n$’s. That is indicated in Fig. 6 where the boundary of the probability distribution spreading, pointed out by the sharp decrease of this quantity, does not correspond to the boundary of the second cell in the case $\delta \omega = 0$, marked by the second arrow on the figure. As follows from calculations with decreasing $V_0$ (or increasing $\delta \omega$) the boundary of the probability distribution spreading decreases. The same is characteristic for the classical accidentally degenerate system [8] where in the case $\delta \omega \neq 0$ the resonance occurs only in several resonance cells in the phase space and the resonances are separated from each other by invariant curves. According to equation (8), there is only one resonance at chosen parameters, and the probability distribution does not spread via the cells, no matter how long we observe the dynamics. This is natural because, as was shown above, the wave with a small amplitude $V_0$ does not affect the states with large Landau numbers $n$. That can be also understood from the expression for the Green function (11). As one can see from Eq. (4), all the QE eigenstates with Landau numbers $n$ from the region $n\hbar \delta \omega \gg V_0$ are localized, $A^q_n = \delta_{n,q}$, with the characteristic localization length $\Delta n \sim V_0/\hbar \delta \omega$ (that can be shown for example in the quasiclassical limit). On the other hand in order for any state to evolve from level $n$ to level $n'$ both $A^q_n$ and $A^q_{n'}$ must be large, but if $A^q_n = \delta_{n,q}$, then $A^q_{n'} = 0$ for all $n' \neq n$. Hence boundary of the probability distribution spreading should be determined from the condition of localization of the QE eigenfunctions (the condition $n\hbar \delta \omega \geq V_0$ was satisfied for the boundary of the probability distribution spreading in Fig. 6). The steep decay of
the probability distribution in Fig. 6 in the quasiclassical approximation corresponds to absence of tunneling between invariant curves. A qualitative comparison of permeabilities of the invariant curves in the case \( \delta \omega \neq 0 \) with permeabilities of the "quantum separatrices" in the case \( \delta \omega = 0 \) (see Fig. 2) indicates that the latter are much more transparent (compare also with Ref. [1]).

In conclusion we would like to do the following remark. The nonperturbed Hamiltonian \( \hat{H}_0 \) yields an unlimited (from above) set of Landau levels with equal distance \( \hbar \omega_c \) between them. When we add a monochromatic wave with the same frequency \( \omega = \omega_c \), a particle can pass to other levels and an interesting question arises: is the diffusion limited or not? The answer depends on how the perturbation depends on coordinate \( x \). If the dependence is linear then quantum states are delocalized and the particle goes to infinity in the limit \( t \to \infty \) (as in the classical case) [10, 11]. As follows from the results of our work, thanks to nonlinearity of the potential \( V_0 \cos(kx - \omega t) \) the quantum states (as classical ones) turn out to be localized. However, the peculiarity of localization in the latter case is that the sharp decay of the probability distribution occurs not everywhere, but only at the boundaries of the quantum resonance cells.

In the next paper we will focus on studying destruction of the quantum resonances under the influence of nonresonant terms in the Hamiltonian (1) when the amplitude of perturbation is not small.

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FIG. 1. Contour plots in the coherent state representation for the eigenstates with eigenvalues at the top and bottom of the QE spectrum $N = 100$, $h = 0.52$ (a). Poincare surfaces of section for one cell for the same parameters of the Hamiltonian (b).

FIG. 2. Probability distribution for seven quantum cells, $h = 0.6$, $V_0 = 0.1$, the number of levels $N = 432$, (a) $t = 4 \times 10^5$, (b) $t = 7 \times 10^5$. The initial state was placed in the center of the first cell.

FIG. 3. The average probability distribution for (a) the first cell $P_1(t)$ and (b) the second cell $P_2(t)$. The parameters and the initial conditions are the same as in Fig. 2.

FIG. 4. Time-averaged probability distribution for two values of the effective Plank's constant $h$ versus cell number. The parameters are the same as in Figs. 2, 3.

FIG. 5. Oscillations of the penetration coefficient from second to the subsequent cells, $t = 4 \times 10^4$, $N = 100$ (three cells), $V_0 = 0.1$. 

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FIG. 6. Probability distribution for the near resonance case \( \delta \omega = 0.003 \), \( t = 4 \times 10^6 \), \( h = 0.52 \), \( V_0 = 0.1 \), the number of levels \( N = 100 \). The initial state in the form \( C_n(0) = \delta_{n,n_0} \) was situated at the level \( n_0 = 6 \).
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