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Exceptional points in bichromatic Wannier–Stark systems

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Received 8 June 2011, in final form 5 September 2011
Published 25 October 2011
Online at stacks.iop.org/JPhysB/44/225301

Abstract
The resonance spectrum of a tilted periodic quantum system for a bichromatic periodic potential is investigated. For such a bichromatic Wannier–Stark system, exceptional points, degeneracies of the spectrum, can be localized in parameter space by means of an efficient method for computing resonances. Berry phases and Petermann factors are analysed. Finally, the influence of a nonlinearity of the Gross–Pitaevskii type on the resonance crossing scenario is briefly discussed.

1. Introduction
The physics of ultracold atoms and Bose–Einstein condensates in optical lattices has made an enormous progress in the last decade. Due to the possibility of controlling all experimental parameters accurately over wide ranges and monitoring the dynamics of the atoms \textit{in situ}, optical lattices have become one of the most prominent model systems in quantum optics, solid state physics and nonlinear dynamics. Nowadays, the shape of the optical potential can be engineered with astonishing precision, including in particular bichromatic optical lattices [1–3]. The manipulation of matter waves in the lattice is routinely accomplished by a static or time-dependent external field, either in an accelerated horizontal lattice [1, 2, 4, 5] or a vertical lattice subject to gravity [6–9] also supported by magnetic levitation [7, 8]. A weak static field accelerates the atoms up to the edge of the Brillouin zone, where they are reflected leading to a periodic motion called Bloch oscillation. A strong field introduces decay by repeated Landau–Zener tunnelling to higher bands. Thus, optical lattices also became an important model system for the study of decay in open quantum systems [1, 2, 5, 6].

The dynamics of a quantum particle in a tilted periodic structure has been a subject of intensive theoretical investigations starting from Bloch’s seminal paper on electrons in crystals [10]. The basic structure of the spectrum and dynamics was then clarified in a discussion by Zak and Wannier [11–13]: in particular, the spectrum is continuous with embedded resonance eigenstates—the so-called Wannier–Stark resonances. These eigenstates are arranged in ladders, where the climbing of the ladder is realized by a translation over one lattice period. Each ladder can be roughly associated with a Bloch band in the field-free case. The equidistant spacing of energies in one ladder leads to a fully periodic motion—the celebrated Bloch oscillations. A review of these results as well as a surprisingly effective algorithm to calculate Wannier–Stark resonance states is given in [14].

In the present paper, we investigate a rather peculiar feature of tilted bichromatic lattices, the existence and properties of exceptional points (EPs). In an open system, the eigenenergies become complex valued, where the imaginary part gives the decay rate (for a comprehensive discussion of such non-Hermitian quantum systems, see the recent textbook by Moiseyev [15]). If a system parameter is varied, as for example the external field strength, the eigenvalues show avoided crossings. For resonance states, two types of level crossing scenarios exist—either the real parts of the energies anti-cross and the imaginary parts cross, denoted as a type I crossing, or vice versa, a type II crossing [16]. A full coincidence of the real and imaginary parts is possible at isolated points in parameter space, the so-called exceptional points. The existence of these points has a remarkable implication on the dynamics of the system. Suppose the
system parameters are varied adiabatically along a cyclic path around an EP. In general, a cyclic evolution leaves the quantum state invariant up to a geometric phase, or Berry phase, which is of great importance both from a fundamental viewpoint as well as for applications [17, 18]. Cycling around an EP has an even stronger effect: in addition to a geometric phase, it exchanges the two crossing states.

Both crossings of type I and II appear for Wannier–Stark systems where the periodic potential is sinusoidal (see, e.g., [14]), but EPs have not been detected. A modulated periodic potential with two additional parameters is, however, flexible enough to allow for such fascinating degeneracies and the first results have been reported recently for a bichromatic potential [19]. Such systems may be very well suited for experimental studies of the phenomena generated by such degeneracies for quantum systems, where experiments are still rare.

In the following, we will discuss the energy spectrum of a quantum particle in a tilted bichromatic optical lattice in the dependence of the system parameters. A pronounced feature of such a Wannier–Stark system is the enhancement of the decay rate by resonant tunnelling (RET) when the feature of such a Wannier–Stark system is the enhancement of the decay rate by resonant tunnelling (RET) when the characteristic width of the discrete distribution of the Fermi level is given by the grid period. In addition, the decay rate can be resonantly enhanced at crossings of the energy surfaces are discussed in detail. EPs are most clearly identified by the divergence of the embedding because experiments are still rare.

2. Wannier–Stark resonances

As mentioned in the introduction, we will demonstrate the existence of EPs for a tilted periodic potential with modulated potential minima and maxima, a bichromatic potential. However, before we turn to the bichromatic case, we first review some general properties of the Wannier–Stark Hamiltonian

\[ H = -\frac{\hbar^2}{2m} \frac{d^2}{dx^2} + V(x) + Fx, \]

where the potential \( V(x + d) = V(x) \) has a single period \( d \). It was shown [11–13, 21] that Hamiltonian (1) with \( F > 0 \) has a continuous spectrum in which discrete ladders

\[ \varepsilon_{a,n} = E_{a,0} + n dF - i \Gamma_a/2 \]

of complex-valued resonances are embedded, where \( a = 1, 2, \ldots \) is the ladder index and \( n \) is the site index. The decay rates \( \Gamma_a/\hbar \), which are the same for all resonances in one ladder, are due to the finite probability for tunnelling out of the lattice towards \( x \to -\infty \) where the Stark term \( Fx \) goes to \( -\infty \) (for positive fields, \( F > 0 \)) so that there is no reflection. The resonance energies (2) and corresponding wavefunctions thus satisfy the non-Hermitean eigenvalue problem

\[ H \Psi_{a,n}(x) = \varepsilon_{a,n} \Psi_{a,n}(x) \]

with purely outgoing (Siegert) boundary conditions \( \partial_x \Psi_{a,n}(x) \to -i k(x) \Psi_{a,n}(x) \) for \( x \to -\infty \) with the local wavenumber \( k(x) = \sqrt{2m(\varepsilon_{a,n} - V(x) - Fx)/\hbar} \). At \( x \to \infty \), the wavefunction satisfies the usual bound state condition \( |\Psi_{a,n}(x)|^2 \to 0 \). Equivalent ways of avoiding reflection at \( x \to -\infty \) include the use of complex absorbing potentials (see [15, 19] and section 4) or complex scaling of the coordinate \( x \) [15, 22]. Here, we use an efficient calculation method introduced in [22, 23] based on a finite basis expansion in momentum space.

The Wannier–Stark ladder (2) can be understood by considering the commutator

\[ [H, T_\ell] = -\ell \ h dF T_\ell \]

between the Wannier–Stark Hamiltonian (1) and the translation operator \( T_\ell \) over \( \ell \) lattice sites. Equation (4) expresses the fact that translation by \( \ell \) lattice sites has the same effect as an energy shift \( \ell F \). Thus, we obtain

\[ \partial_x \Psi_{a,n}(x) = T_\ell H \Psi_{a,n}(x) + [H, T_\ell] \Psi_{a,n}(x) \]

\[ = (\delta_{a,a} - \ell dF) \Psi_{a,n}(x), \]

which leads to ladder (2) of resonance energies with the respective wavefunctions

\[ \Psi_{a,n}(x) = \Psi_{a,0}(x - n dF). \]

The general dependence of the decay rates on the field strength is approximately given by \( \Gamma(F) \propto F \exp(-\pi D^2/F) \) where \( D \) is the energy gap between the ground and first excited Bloch band. This result is obtained by Landau–Zener theory under the assumption that decay is mainly determined by tunnelling from the ground band to the first excited band as successive tunnelling events into higher bands, which finally lead to decay towards \( x \to -\infty \), are fast compared to the first tunnelling process [24, 25]. Deviations from the Landau–Zener dependence occur if a state of a lower ladder with energy \( \varepsilon_{a,n} \) is in resonance with a state of a higher ladder at a different site, i.e. \( E_{a,n} = E_{a',n'} \) [5, 14]. Such a resonant coupling between two ladders leads to a strong increase of the decay rate of the lower ladder so that it is called resonantly enhanced tunnelling (RET). At the same time, the decay rate in the upper ladder is lowered. Whenever two ladders are in resonance, the decay rate for the lower ladder shows a peak, whereas the upper band has a dip.

Now we turn to the bichromatic or double-periodic potential

\[ V(x) = \frac{V_0}{2}(\cos(2\pi x/d) + \delta \cos(4\pi x/d + \phi)), \]

given by the sum of a grid with period \( d \) and an additional \( d/2 \)-periodic grid which creates \( d \)-periodic potentials with modulated minima and maxima of different heights, depending on \( \phi \). Throughout this paper, we use scaled units such that \( d \) is equal to \( 2\pi \) and \( m = \hbar = 1 \). The energies are measured.
in units of $8E_R$, where $E_R = \hbar^2\pi^2/(2md^2)$ is the recoil energy. In the subsequent computations, we furthermore fix the potential strength as $V_0 = 1$.

The band structure of the double-periodic potential (7) is characterized by the splitting of the ground band into two minibands. The Wannier–Stark ladder splits into two miniladders, where the width of the splitting depends on the modulation $\delta$. In the limit of vanishing modulation $\delta \to 0$, the miniladders are energetically degenerate. In analogy to equation (4), one can derive the relations

$$[H, T_{2\ell}] = -(2\ell + 1) dF T_{\ell}$$

(8)

$$[H, T_{2\ell+1}G] = -(2\ell + 1) dF T_{\ell}$$

(9)

(see [26, 27] for details), where $H$ is the Hamiltonian, $T_{\ell}$ is the translation operator over $\ell$ lattice sites and $G$ is an operator that switches the sign of the modulation $\delta$ in all following terms. Using (8) and (9), one can show that the Wannier–Stark ladder of eigenenergies of the unperturbed system splits into the two miniladders

$$\mathcal{E}_{\alpha,2\ell} = \mathcal{E}_\alpha(\delta) + 2\ell dF$$

(10)

$$\mathcal{E}_{\alpha,2\ell+1} = -\mathcal{E}_\alpha(\delta) + (2\ell + 1) dF.$$  

(11)

The energy offset $\mathcal{E}_\alpha(\delta)$ is an antisymmetric function of $\delta$. Each Wannier–Stark ladder bifurcates into two which are energetically shifted with a distance $2\mathcal{E}_\alpha(\delta)$.

Let us first consider the topology of the energy surfaces in an arbitrarily chosen part of the parameter space. Figure 1 shows the imaginary part of the eigenenergies, i.e. the decay rate of the two most stable states in dependence on the two parameters $\phi$ and $1/F$ where the modulation $\delta = 1$ is kept fixed.

Figure 1. Decay rate $\Gamma_{1/2}$ as a function of the inverse field strength $1/F$ and the phase $\phi$ with $\delta = 1$. The configurations marked by a vertical line correspond to EPs.

If we vary only $1/F$ keeping the potential, i.e. $\phi$, fixed, we observe the familiar RET spectra with eigenvalue crossings and avoided crossings. For the cut with $\phi = -1$ appearing at the front of the surfaces in figure 1, for example, we have a series of avoided crossings of the imaginary parts, whereas the real parts cross, i.e. a type II crossing scenario. Cuts at different values of $\phi$ show type I crossings. One can straightforwardly verify that the two eigenfunctions involved are also degenerate for these parameters so that points (12) and (13) are indeed EPs. This is illustrated in figure 2 which exemplarily shows the two Wannier–Stark eigenfunctions $\Psi_1$ and $\Psi_2$ for parameters (12).

In the following, we address the question how the emergence of EPs in a bichromatic Wannier–Stark potential can be understood in terms of the shape of the potential.

3. Exceptional points

3.1. Overview

EPs occur in systems described by non-Hermitian Hamiltonians depending on a set of parameters. As mentioned above, a point in parameter space at which both the complex eigenvalues and the eigenstates of two different eigenmodes coincide is called an EP [15, 28, 29]. EPs must be distinguished from so-called diabolic points that occur if there is a coincidence of the eigenvalues but not of the eigenstates. Diabolic points are more familiar from Hermitian systems, but in fact they are more special than the EPs and can be viewed as a coincidence of two EPs (see, e.g., [16, 30] and references therein). Bichromatic optical lattices showing diabolic points in the band structure have been proposed as quantum simulators for the Dirac equation [31].

Before we introduce a systematic algorithm for finding the EPs for the tilted periodic potential (7), we first have a look at the set of solutions $\mathcal{E}(1/F, \delta, \phi)$ shown in figure 1. EPs are searched by means of the criterion that the energy difference $\Delta\mathcal{E} = |\mathcal{E}_1 - \mathcal{E}_2|$ between the two lowest levels must vanish. This criterion is satisfied for the following points marked by vertical lines in the figure:

$$1/F = 3.769, \quad \delta = 1, \quad \phi = -2.991.$$  

(12)

$$1/F = 6.662, \quad \delta = 1, \quad \phi = -2.228.$$  

(13)

One can straightforwardly verify that the two eigenfunctions involved are also degenerate for these parameters so that points (12) and (13) are indeed EPs. This is illustrated in figure 2 which exemplarily shows the two Wannier–Stark eigenfunctions $\Psi_1$ and $\Psi_2$ for parameters (12).

3.2. Emergence of EPs in tilted optical lattices

In the following, we address the question how the emergence of EPs in a bichromatic Wannier–Stark potential can be understood in terms of the shape of the potential.
The definition of an EP requires the coincidence of both the energies and eigenfunctions of the two states considered.

Let us first recall the features of the single-periodic Wannier–Stark system introduced in section 2, where a simultaneous degeneracy of the eigenenergies and eigenfunctions is impossible. The Stark potential $Fx$ leads to a localization of the eigenfunctions and induces a ladder structure of the eigenenergies $E_{\alpha,n} = E_{\alpha,0} + 2\pi F n$ (see equation (2)). For a finite field $F > 0$, the energies within a ladder differ by multiples of $2\pi F$, such that a degeneracy in $E$ can only occur between states of different ladders, resulting in resonantly enhanced tunnelling. However, the states involved are localized in different wells of the potential, such that they cannot be described by the same wavefunction.

The situation is different in a bichromatic lattice. Within one period of the potential $V(x)$, there are two potential minima, in which a Wannier–Stark resonance is localized. By varying one parameter, e.g. the field strength $F$, one can achieve a degeneracy of the eigenenergies. Now it is possible to also fine-tune the remaining parameters to realize a full coincidence as illustrated in figure 4. An EP is found when the more stable state is destabilized and vice versa while keeping their energies degenerate.

At the EP, both the eigenenergies and the eigenstates must coincide. Therefore, the overlap integral

$$S_{1,2} = |\langle \Psi_1 | \Psi_2 \rangle| \leq 1$$

of the two (normalized) eigenstates should approach unity at the EP. For parameters far away from it, the overlap is expected to be small. As an example, figure 3 displays the overlap in dependence on $1/F$ for system parameters at some distance from an EP (left panel), where we find a rather small peak of the overlap. The right panel of figure 3 displays the overlap $S$ in dependence on $F$ in the vicinity of EP (12). For $1/F \approx 3.7$, which corresponds to the configuration of the EP, we observe a pronounced peak with $S \approx 1$ as the two wavefunctions are degenerate. This behaviour can be used as a tool for detecting an EP. Note that alternatively one might consider the Petermann factor $K_\alpha$ describing the ‘(self-)overlap’ between left and right eigenvectors that we will use in section 4.

To further investigate this phenomenon, figure 4 shows both the potential and the wavefunction for the parameters

$$1/F = 3.814, \; \delta = 2.251, \; \phi = -3.035$$

that correspond to an EP as well as for a slightly different parameter set in the vicinity of the EP (see table 1 for the exact configurations). In agreement with our previous observations, the overlap only assumes values of $S \approx 1$ for a very narrow
interval of field strengths $F$. The potentials shown in the two panels in figure 4 are hardly distinguishable at first glance. Yet there is a significant difference as far as the corresponding pairs of Wannier–Stark eigenfunctions and in particular their overlap are concerned. The difference is also evident in the corresponding complex eigenenergies given in table 1. Looking at the decay rates, we observe that at the transition to the EP, the ground state becomes less stable, whereas the lifetime of the excited state increases.

The emergence of the degeneracy can be related to the shape of the double-periodic potential. Within a period of the potential $V(x)$, we have two minima marked by $\mathcal{\Omega}$ and $\mathcal{\Omega}'$ in the left lower panel of figure 4. These minima are slightly shifted with respect to each other. For $V_{\text{min}}^{\mathcal{\Omega}} < V_{\text{min}}^{\mathcal{\Omega}'}$ the ground state wavefunction is mainly localized in the first well (cf. the left panel of figure 4). Together with well $\mathcal{\Omega}'$ one obtains the shape of a double well potential. The first excited state reveals a high probability density in the second well. Since the barrier in the direction $x \to -\infty$ is lower for well $\mathcal{\Omega}$ the probability for tunnelling processes in this direction is higher leading to a larger decay rate (cf. table 1).

For the field strength $1/F \approx 3.814$ corresponding to the configuration of the EP, the ground and first excited resonance states are energetically fully degenerate. The first excited state is stabilized since tunnelling towards the ‘more stable well’ $\mathcal{\Omega}$ is enhanced. At the same time, the ground state is destabilized because the probability for tunnelling from well $\mathcal{\Omega}'$ towards $x \to -\infty$ is likewise enhanced. The EP corresponds to the balanced situation where the two Wannier–Stark states become indistinguishable so that there is a degeneracy between wavefunctions belonging to different bands which cannot be achieved in the single-periodic Wannier–Stark system.

### Table 1. System parameter sets in the vicinity of EP (15).

<table>
<thead>
<tr>
<th>$1/F$</th>
<th>$\delta$</th>
<th>$\phi$</th>
<th>$E_1$</th>
<th>$\Gamma_1$</th>
<th>$E_2$</th>
<th>$\Gamma_2$</th>
</tr>
</thead>
<tbody>
<tr>
<td>3.000</td>
<td>2.251</td>
<td>-3.141</td>
<td>$2.561 \times 10^{-1}$</td>
<td>$2.961 \times 10^{-2}$</td>
<td>$1.858 \times 10^{-3}$</td>
<td>$1.361 \times 10^{-1}$</td>
</tr>
<tr>
<td>3.814</td>
<td>2.251</td>
<td>-3.035</td>
<td>$1.538 \times 10^{-1}$</td>
<td>$7.427 \times 10^{-2}$</td>
<td>$1.538 \times 10^{-1}$</td>
<td>$7.427 \times 10^{-2}$</td>
</tr>
</tbody>
</table>

3.3. Systematic search for EPs

So far all EPs have been determined in the same way: if one suspects the existence of an EP in some part of the configuration space, eigenvalues and eigenfunctions of the Hamiltonian are analysed with respect to the conditions

| $|E_1 - E_2| = 0$ | $(16)$ |
| $S_{1,2} = |\langle \Psi_1 | \Psi_2 \rangle| = 1.$ | $(17)$ |

In the present study, we numerically minimized the eigenvalues distance in (16) using a standard method based on the simplex algorithm (see, e.g. [32, 33]). Finally it was checked if the overlap condition (17) is satisfied as well. Note that a more efficient method for localizing EPs has been developed recently [34, 35].

So far, one of the three parameters ($1/F, \delta, \phi$) has been kept fixed, whereas the other two were varied. Now we search for EPs in the fully three-dimensional parameter space by solving the minimization problem (16). In order to converge to the desired EP instead of some local minimum, the algorithm requires a suitable initial guess for the parameters ($1/F, \delta, \phi$). To this end, we represent the parameter space in terms of spherical coordinates

$$\Delta(1/F) = r \sin \vartheta \cos \varphi$$
$$\Delta \delta = r \sin \vartheta \sin \varphi$$
$$\Delta \phi = r \cos \vartheta$$

(18)

If the parameters $r_{\text{EP}} = (1/F_{\text{EP}}, \delta_{\text{EP}}, \phi_{\text{EP}})$ of an EP are known, one can look for further solutions $r_{\text{EP}} = \text{EP} + \Delta r$ in its vicinity. Thus, we search for EPs in the following manner:

(i) find an EP (e.g. with the methods described further above);
(ii) choose some small ‘distance’ r, in which to look for the next EP (see below);
(iii) minimize condition (16) using the simplex algorithm with the parametrization ($r = \text{const.}, \vartheta, \varphi$);
(iv) repeat this procedure starting with the configuration of the newly found EP.

For the sake of completeness, we note that we use

$$r(\Delta(1/F), \Delta \delta, \Delta \phi) = \sqrt{(\Delta(1/F))^2 + (\Delta \delta)^2 + (\Delta \phi)^2}$$

(19)

with $\Delta(1/F) = 1/F_1 - 1/F_2$, $\Delta \delta = \delta_1 - \delta_2$ and $\Delta \phi = \phi_1 - \phi_2$ as the distance between two configurations ($1/F_1, \delta_1, \phi_1$) and ($1/F_2, \delta_2, \phi_2$).

With this procedure the minimization problem (16) is solved in a two-dimensional subspace (the surface of a sphere) of the three-dimensional parameter space which reduces the numerical effort. The finite area of that subspace also improves the convergence. This systematic search detects all EPs between the two ladders in the parameter region under consideration. Obviously, further EPs can exist between higher ladders, which are not detected. However, these are of less practical interest, as the decay is significantly stronger.

Starting with EP (12), we obtain a one-dimensional manifold formed by the EPs of the two most stable resonances of the double-periodic Wannier–Stark system that is embedded in the three-dimensional parameter space. This is displayed in figure 5. In addition, we have included a contour plot of the absolute difference $\Delta \Gamma = |\Gamma_1 - \Gamma_2|$ of the decay rates on the plane $\delta = 1$. The intersections of the EP curve with this plane are located at the zeros of (16). For $\delta = 1$, three solutions can be detected in the region considered.

One can draw the following conclusions: in the single-periodic system, no EPs were found. The results for the bichromatic system reveal the transition between the two different potentials: in the limit $\delta \to 0$, the position of the EP is shifted towards an infinitely strong field strength $F (1/F \to 0)$, so that an EP cannot exist in the single-periodic limit.

Finally we briefly discuss the shape of the curve. From figure 5 we see that the EP curve bends strongly for $\delta \approx 2.3$, so
that no EPs can be found for higher values of the modulation $\delta$. Yet for small field strengths $F$, i.e. for high values of $1/F$, there are intersections of the EP curve with the planes $(1/F, \delta = \text{const}, \phi)$. This local behaviour repeats itself and can also be observed for higher bands. The results in the domain of interest support our expectation that there are no EPs in the limit $F \to 0$, i.e. without an external field. To gain further insight into the characteristic shape of the EP curve, additional investigations are needed.

### 3.4. Cyclic parameter variation

In the following we study the behaviour of eigenvalues and eigenvectors under cyclic parameter variation. As an example, we consider the vicinity of the EP (12) and choose $\delta = 1$ to be constant. The quantities $F$ and $\phi$ are varied along the paths

$$
\frac{1}{F} = \frac{1}{F_\text{EP}} + r \sin(\beta) + \frac{1}{F_0}
$$

$$
\phi = \phi_\text{EP} + r \cos(\beta) + \phi_0
$$

(20)

parametrized by an angle $\beta \in [0, 2\pi]$. As in subsection 3.3, $r$ is a numerical quantity that can be identified as a radius in parameter space.

Apart from paths that enclose the EP $(1/F_\text{EP}, \delta_\text{EP}, \phi_\text{EP})$, we are also interested in paths that do not enclose it, which can be controlled via the shifts $F_0$ and $\phi_0$. Figure 6 shows nine selected paths; the corresponding eigenvalue trajectories in the complex energy plane are displayed in figure 7.

We observe two qualitatively different kinds of behaviours: any closed path in parameter space that does not enclose the EP leads to eigenvalue trajectories in the complex energy plane that consist of two separated closed curves, one for the ground state and one for the first excited state. Thus, the eigenvalues always remain on the same Riemann sheet. A different result is obtained for path ➀ that encloses the EP. Instead of two separate eigenvalue curves in the complex energy plane, we now find a single closed curve consisting of the complex eigenvalues of both eigenstates involved. This property is typical of EPs [36] and provides a useful criterion for proving the existence of an EP in experiments (see e.g. [37, 38]). After two full cycles, the initial eigenvalue is recovered, i.e. the eigenvalues vary with a period of $4\pi$ in parameter space.

It is also instructive to look at the eight outer subplots in figure 7 in the cyclic sequence

$\circled{1} \to \circled{2} \to \circled{3} \to \circled{4} \to \circled{5} \to \circled{6} \to \circled{7} \to \circled{8} \to \circled{1}$. (21)

Comparing the eigenvalue trajectories for path ➀ and path ➀ we observe that the assignment of the two trajectories to the ground state (blue) and the first excited state (red) is interchanged. This can be interpreted as a closed path in parameter space that also encloses the EP and hence moves from one Riemann sheet to the other.

Apart from the eigenvalues we are also interested in the eigenfunctions. In order to monitor their behaviour, it is sufficient to consider only one complex number, namely the projection of each wavefunction on one of the basis vectors in the plane-wave basis chosen for the calculations. Here we choose the projection $|\Psi_0|_{\text{max}}$ with the maximum value of $|\Psi_0|^2$; the other projections show a qualitatively similar behaviour. In figure 8, we display the phase $\Omega$ (left panel) as well as the real and imaginary parts (right panel) of $|\Psi_0|_{\text{max}} = |(\Psi_0)_{\text{max}}| e^{-i\Omega}$ for the ground and first excited state for a cyclic parameter variation around the EP, i.e. when...
Figure 7. Eigenvalue trajectories in the complex energy plane \( \hat{\mathcal{E}} = E - i\Gamma/2 \) for a cyclic parameter variation of \( F \) and \( \phi \) for \( \delta = 1 \), cf also figure 6. The energy of the EP is marked by \( \otimes \).

Figure 8. Component \( (\Psi_\alpha)_{\text{max}} = |(\Psi_\alpha)_{\text{max}}| e^{-i2\pi\Omega} \) of the ground state (blue) and the excited state (red) for a cyclic parameter variation around EP (12). The phase \( \Omega \) (left panel) as well as the real and imaginary parts (right panel).

The phase \( \beta \) in parameter space varies from 0 to \( 2\pi \). The figures reveal the expected behaviour: after one complete cycle around the EP, the phases \( \Omega \) of the two components are interchanged, however, with a phase change of \( \pi \), i.e. a different sign. In the right panel, showing the eigenvector components in the complex plane, one observes the resulting \( 8\pi \)-periodicity. Starting at the points \( a \) (ground state) and \( d \) (first excited state) for \( \pi = 0 \), the selected components \( (\Psi_1)_i \), reach the \( b \) and \( a \) points after one cycle, where \( b \) corresponds to the sign-changed component \( (\Psi_2)_i \) at \( d \). The dashed lines show the development for three further cycles.

The periodicity of the eigenvalues and eigenfunctions can be summarized in the following diagrams (cf [16]):

\[
\begin{align*}
\left\{ \mathcal{E}_1, \mathcal{E}_2 \right\} & \xrightarrow{4\pi} \left\{ \mathcal{E}_2, \mathcal{E}_1 \right\}, \\
\left\{ \Psi_1, \Psi_2 \right\} & \xrightarrow{8\pi} \left\{ -\Psi_2, -\Psi_1 \right\}.
\end{align*}
\]
where $\Psi_1$ and $\Psi_2$ are the Wannier–Stark eigenstates involved in the formation of the EP, and $\mathcal{E}_1$ and $\mathcal{E}_2$ are the corresponding eigenenergies.

Finally, looking again at figure 5, we observe that in the vicinity of the parameter values $1/F = 4$, $\delta = 2.3$ and $\phi = -3$, two EPs approach each other. This region is magnified for $\delta = 2.3$ in figure 9, which shows the energy difference $\Delta \mathcal{E} = |\mathcal{E}_1 - \mathcal{E}_2|$ in dependence on $1/F$ and $\phi$. It is instructive to have a closer look at the surfaces of the real and imaginary parts of the eigenenergies in parameter space in this region. This is illustrated in figure 10, which shows the Riemann surfaces of the complex energy in the vicinity of these two EPs. The observed local behaviour resembles the ‘double-coffee-filter’ scenario typical of unfolding a diabolic point into two EPs [16, 30, 39–41]. Note that the apparent gaps in the Riemann surfaces are an artefact of the finite discretization of the parameter space.

### 3.5. The Petermann factor

For a non-Hermitian system, one has to distinguish between right and left Wannier–Stark states, i.e. the eigenstates of $H$ for the eigenvalue $E_\alpha$ and of $H^\dagger$ for the eigenvalue $E^*_\alpha$, denoted as $|\Psi_\alpha\rangle$ and $|\bar{\Psi}_\alpha\rangle$, respectively. These states form a bi-orthogonal set, i.e. $\langle \bar{\Psi}_\alpha | \Psi_\alpha \rangle = 0$ for $\alpha \neq \alpha'$. At an EP, two eigenstates of $H$ coincide, as well as the corresponding eigenstates of $H^\dagger$, with the consequence that the scalar product $\langle \bar{\Psi}_\alpha | \Psi_\alpha \rangle$ vanishes. A useful quantity in this context is the Petermann factor [42, 43]

$$K_\alpha = \frac{\langle \bar{\Psi}_\alpha | \bar{\Psi}_\alpha \rangle \langle \bar{\Psi}_\alpha | \Psi_\alpha \rangle}{|\langle \bar{\Psi}_\alpha | \Psi_\alpha \rangle|^2},$$

which provides a measure of the overlap between left and right eigenvectors. Note that $K_\alpha$ is independent of the normalization of the states and diverges at the EP. The Petermann factor provides a convenient tool for localizing an EP.

Figure 11 shows the Petermann factor of the ground state and the first excited state in dependence on the field strength $F$. The modulation $\delta$ and phase $\phi$ were chosen according to the configuration of EP (12). The offset between the two curves is caused by the difference in the decay rates of the two respective states. At the position corresponding to the EP, both Petermann factors $K_1, K_2$ show pronounced peaks. Due to the discretization of the field strength and the fact that configuration (12) is only very close to but not exactly in the EP, the numerical value of $K_\alpha$ remains, of course, finite.
4. Nonlinear crossing scenarios

One of the most promising realizations of Wannier–Stark systems, both theoretically and experimentally, is based on atomic Bose–Einstein condensates in tilted optical lattices. Near the zero-temperature limit, these systems can be described by a NLSE or Gross–Pitaevskii equation (GPE) (see e.g. [44–47])

\[
\left(-\frac{\hbar^2}{2m}\frac{\partial^2}{\partial x^2} + V(x) + Fx + g|\Psi(x)|^2\right)\Psi(x) = \mu \Psi(x)
\]

(25)

in a mean-field approach, where the nonlinear term \(g|\Psi(x)|^2\) takes into account the interaction between the condensate particles. The interaction between the particles can either be repulsive \((g > 0)\) or attractive \((g < 0)\). In the case of Wannier–Stark resonances, the chemical potential \(\mu = M - i\Gamma/2\) is complex where the real part \(M\) describes the position of the resonance and the imaginary part \(\Gamma/2\) accounts for the decay rate.

To date, complex eigenvalue crossing scenarios and EPs in the context of nonlinear systems have mostly been investigated by means of simple nonlinear non-Hermitian two-level models [48–50], three-level systems [51], or to some extent analytically solvable model systems such as one-dimensional \(\delta\) and \(\delta\)-shell potentials [52] or multiple-barriers [53] of a delta-comb [54]. An exception is the discovery of EPs for atomic gases with an attractive \((g < 0)\) or repulsive \((g > 0)\) nonlinearity in opposite directions for positive or negative values of \(g\), assuming that the existence of the EP is not affected by weak interactions. Figure 13 displays the Petermann factors of the ground state \(K_1\) and of the first excited state \(K_2\) for a parameter variation of \(F\) around EP (12) of the linear system. The interaction strength was chosen as \(g = -0.02\) (blue), 0 (black) and 0.02 (red).

The numerical calculation of nonlinear Wannier–Stark resonances is a nontrivial task because most methods established for the linear system cannot be straightforwardly adapted since the superposition principle is no longer valid. In [55, 56], nonlinear Wannier–Stark resonances of the ground ladder in single-periodic and double-periodic systems were calculated using a complex scaling procedure developed in [57, 58]. Here we use a method based on complex absorbing potentials in position space [19] as it is the only procedure so far that has successfully been used to calculate Wannier–Stark resonances of excited ladders. In the following we briefly discuss the impact of the nonlinear interaction on eigenvalue crossing scenarios and EPs for the bichromatic WS system discussed in the preceding sections.

As a first example, figure 12 shows the real and imaginary parts of the eigenvalues of the chemical potential \(\mu\) in dependence on the static field strength \(F\) for a WS system again in the vicinity of EP (12). The black curves represent the real and imaginary parts of the resonance energies of the two most stable states for the non-interacting system \((g = 0)\) which cross at \(F = 0.2653\). A moderate nonlinearity removes this degeneracy and we observe a type I crossing (cf section 1) for a repulsive interaction \((g = +0.02,\) red curves\)) and a type II crossing for an attractive interaction \((g = -0.02,\) blue curves\)) in agreement with the results already briefly reported in [19].

This indicates that the position of the EP is shifted by the nonlinearity in opposite directions for positive or negative values of \(g\), assuming that the existence of the EP is not affected by weak interactions. Figure 13 displays the Petermann factors \(K_1\) and \(K_2\) for the resonance states shown in figure 12. The Petermann peaks for \(g = \pm 0.02\) are now strongly reduced and shifted to stronger fields for repulsive and weaker fields for attractive interactions. Moreover, the two maxima are also shifted relative to each other: the maxima of the excited state are more affected by the nonlinearity than the ground state.

Finally we have a closer look at the range of moderately increased attractive interaction strengths \(g\) with \(-0.3 < g < 0\). The corresponding results are shown in figure 14 for the ground
interaction strength is chosen as $g_F$ around the configuration of EP (12) in the linear case. The exists for all values of $g_f$ familiar behaviour; they form a single closed curve for paths space were considered. The eigenvalue trajectories show the trajectories in the complex plane for closed paths in parameter space coincide as they localize in different lattice sites. Furthermore, can only occur for eigenstates of different ladders which cannot eigenfunctions. It was demonstrated that in the limit case as well as the properties of the corresponding degenerate location within the parameter space was discussed in detail had so far mostly been considered for very simple toy models. This paper investigated the spectrum of an experimentally realizable model system for decay in open quantum systems, namely a quantum particle in a tilted bichromatic lattice. The analysis concentrated on the EPs of the system, i.e. points in parameter space where both the complex eigenvalues and the convergence properties of the method for calculating nonlinear Wannier–Stark resonances. Thus, substantial modifications of the presented numerical procedures are required. In this context, the criterion of a vanishing inverse Petermann factor $1/K_\alpha$ for the configuration of an EP might help. One advantage of this criterion is that it only requires the eigenstates of a single ladder.

5. Conclusion

This paper investigated the spectrum of an experimentally realizable model system for decay in open quantum systems, namely a quantum particle in a tilted bichromatic lattice. The analysis concentrated on the EPs of the system, i.e. points in parameter space where both the complex eigenvalues and the eigenfunctions of two different eigenmodes coincide, which had so far mostly been considered for very simple toy models. An efficient method for finding EPs was presented and their location within the parameter space was discussed in detail as well as the properties of the corresponding degenerate eigenfunctions. It was demonstrated that in the limit case of a monochromatic lattice, there can be no EPs for finite values of the tilt, because degeneracies of the eigenenergies can only occur for eigenstates of different ladders which cannot coincide as they localize in different lattice sites. Furthermore, the geometric phases (Berry phases) and the eigenvalue trajectories in the complex plane for closed paths in parameter space were considered. The eigenvalue trajectories show the familiar behaviour; they form a single closed curve for paths in parameter space that enclose an EP but two distinct closed curves otherwise. In particular, the two crossing states are interchanged after a full cycle in parameter space enclosing an EP. Finally the case of an interacting Bose–Einstein condensate in a tilted bichromatic optical lattice was discussed within a mean-field approach by solving the corresponding NLSE or GPE. It was found that the type of crossing, i.e. if the real parts of the eigenenergy anti-cross while the imaginary parts (decay rates) cross (type I) or vice versa (type II), depends on the sign of the interaction between the particles (i.e. whether it is attractive or repulsive). These crossing scenarios, as well as an analysis of the Petermann factor, which measures the overlap between left and right eigenstates and shows a pronounced peak in the vicinity of an EP, indicate that the EPs are shifted by the nonlinear interaction in a direction opposite to its sign.

Acknowledgments

We would like to thank Sandro Wimberger for fruitful discussions. Financial support by the Deutsche Forschungsgemeinschaft (DFG) via the Graduiertenkolleg 792 ‘Nichtlineare Optik und Ultrakurzzeitphysik’ is gratefully acknowledged. KR acknowledges support from a scholarship of the Université Libre de Bruxelles.

References

[10] Bloch F 1928 Z. Phys. 52 555
[18] Sj¨oqvist E 2008 Physics 1 35
[47] Leggett A J 2001 Rev. Mod. Phys. 73 307