Population inversion and dynamical phase transitions in a driven superconductor

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We consider a superconductor in which the density of states at the Fermi level or the pairing interaction is driven periodically with a frequency larger than the superconducting gap in the collisionless regime. We show by numerical and analytical computations that a subset of quasiparticle excitations enter into resonance and perform synchronous Rabi oscillations leading to cyclic population inversion with a frequency that depends on the amplitude of the drive. As a consequence a different “Rabi-Higgs” mode emerges. Turning off the drive at different times and modulating the strength allows access to all known dynamical phases of the order parameter: persistent oscillations, oscillations with damping and overdamped dynamics. We discuss physical realizations of the drive and methods to detect the dynamics.

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I. INTRODUCTION

Quasiparticle relaxation times in superconductors can easily reach nanoseconds at low temperatures [1] while the typical order-parameter dynamics is on the hundreds of femtoseconds time scale [2–6], leaving a large window where, in principle, it is possible to observe energy-conserving out-of-equilibrium dynamics of the superconducting order parameter. Ultracold Fermi gases provide another platform where the order-parameter dynamics of a fermionic condensate can be studied in real time thanks to the possibility to couple directly to [7], or indirectly through the modulation of Feshbach resonances [8,9], and external potentials defined by optical traps [10].

Restricting to Bardeen, Cooper, and Schrieffer (BCS)–like wave functions and within the BCS reduced Hamiltonian, it has been shown [11–13] that there are essentially three zero-temperature dynamical phases, characterized by the qualitatively different time evolution of the order parameter after a sudden change of the interaction. For small perturbations of the BCS ground state, one finds an oscillatory behavior with a frequency $2\Delta_\infty$ and a $t^{-1/2}$ decay of the order parameter to a steady-state value $\Delta_\infty$, smaller than the thermodynamical order parameter $\Delta_0$ at equilibrium. This can be understood appealing to a linearized dynamics where quasiparticle excitations evolve with a frequency determined by its own quasiparticle energy leading to the dephasing of the excitations that build the perturbation and washing out any macroscopic manifestation of the dynamics. On the other hand, if the perturbation is large, two outcomes are possible. If one starts from an initial condition where the order parameter is much larger than the equilibrium value, the dynamics becomes overdamped and the asymptotic stationary value becomes zero. Instead, one starts from an order parameter much smaller than the one at equilibrium, persistent oscillations occur where all quasiparticles evolve synchronously, driven by the pairing field that they build self-consistently [12].

In the case of an isotropic band structure, these spontaneous amplitude modes of the superconducting gap (Higgs-like modes) are totally symmetric, i.e., there is no dependence of the modes on the Euler angles parametrizing a point on the Fermi surface. More importantly, because of the approximate particle-hole symmetry of the BCS state, the charge is locally conserved in each real-space unit cell and therefore long-range effects of the Coulomb interactions are irrelevant and can be neglected from the outset.

In this work we study the effect of a periodic drive [14,15] that couples to the amplitude of the superconducting order parameter. We discuss several type of drives and how they can be realized in practice. In principle, different drives can be implemented in ultracold atoms where the ease of manipulation is one of their well-known characteristics [8,10]. In solid-state systems, on the other hand, one would think that driving requires more effort. However, the first observations of the out-of-equilibrium fermionic condensate dynamics were done in such systems [2,4,5]. Drives that modulate the density of states (DOS) assisted by phonons have been proposed long ago [16,17]. We argue that in unconventional superconductors also phonon assisted drives that modulate the coupling constant are feasible. Other proposed realizations are impulsive stimulated Raman scattering (ISRS) drives [2,3] and THz drives [4–6,18–21]. We find that under such periodic drives (both for ultracold atomic gases and solid-state systems) a subset of quasiparticles, whose energy matches the frequency of the drive, enter into resonance and perform synchronous Rabi oscillations resulting in an oscillation of the amplitude which will be referred to as the “Rabi-Higgs” mode that is clearly distinct from the spontaneous Higgs mode discussed before in the literature [11–13].
An important physical consequence of the Rabi-Higgs mode is that a periodic population inversion is produced for the resonant quasiparticles with the period determined by the amplitude of the drive. This result can be understood mapping the wave function to a system of Anderson pseudospins and applying the rotating wave approximation (RWA) in analogy with Rabi oscillations in nuclear-magnetic-resonance (NMR) experiments. We show that it is possible to explore the full dynamical phase diagram of an out-of-equilibrium fermionic condensate by turning off the drive at different instants of the Rabi cycle, very much as in regular pulsed NMR experiments.

In addition to the neutral spontaneous amplitude modes one can consider also charge modes. The simplest spontaneous charge mode is the longitudinal sound mode which one can consider also charge modes. The simplest spontaneous charge modes are the oscillations observed experimentally with Rabi oscillations in nuclear-magnetic-resonance (NMR) experiments. We show that it is possible to explore the full dynamical phase diagram of an out-of-equilibrium fermionic condensate by turning off the drive at different instants of the Rabi cycle, very much as in regular pulsed NMR experiments.

We will restrict to drives that couple to the pure pseudospin operators are given by

\[ S_k^+ = \frac{1}{2} (c_{k\uparrow} c_{-k\downarrow}^\dagger + c_{-k\uparrow} c_{k\downarrow}^\dagger), \]

\[ S_k^- = \frac{1}{2} (c_{k\uparrow} c_{-k\downarrow}^\dagger - c_{-k\uparrow} c_{k\downarrow}^\dagger), \]

\[ S_k^\sigma = \frac{1}{2} (1 - c_{k\uparrow}^\dagger c_{-k\downarrow}^\dagger - c_{-k\uparrow}^\dagger c_{k\downarrow}^\dagger), \]

where \( c_{k\sigma} \) is the creation (destruction) operator for electrons with momentum \( k \) and spin \( \sigma \).

Due to the infinite range of interactions, assumed in the second term of Eq. (1), the mean-field approximation is exact in the thermodynamic limit and the time-dependent mean-field equations describe the exact dynamics. The BCS mean-field Hamiltonian can be written as

\[ H_{MF} = -\sum_k S_k \cdot b_k, \]

where \( b_k(t) = 2\Delta(t) \hat{x} + 2\Delta'(t) \hat{y} + 2\xi_k(t) \hat{z} \) is an effective magnetic field. The real and imaginary part of the instantaneous superconducting order parameter are defined self-consistently as

\[ \Delta(t) = \lambda(t) S_k^z, \]

\[ \Delta'(t) = \lambda(t) S_k^y, \]

with \( S_k^z = \sum_k \langle S_k^z \rangle \) and \( S_k^y = \sum_k \langle S_k^y \rangle \). Clearly the \( x-y \) projection of the pseudospins is related to the superconducting order parameter while the \( z \) projection is related to charge fluctuations through Eq. (2).

From here on we will denote equilibrium static quantities with a \(^0\) superscript/subscript. Without loss of generality we take \( \langle S_k^y \rangle = 0 \). At equilibrium, in the absence of excitations, the pseudospins align in the direction of their local field \( b_k^0 = 2\Delta_0 \hat{x} + 2\xi_0 \hat{z} \) in order to minimize the system’s energy, which is described by the spin Hamiltonian [Eq. (3)]. The zero-temperature pseudospin texture is given by

\[ \langle S_k^y \rangle^0 = \frac{\Delta_0}{2\sqrt{\xi_k^2 + \Delta_0^2}}, \]

\[ \langle S_k^z \rangle^0 = \frac{\xi_k}{2\sqrt{\xi_k^2 + \Delta_0^2}}. \]

Out of equilibrium the Anderson pseudospins obey the equations of motion for magnetic moments in a time-dependent magnetic field,

\[ \frac{d\langle S_k \rangle}{dt} = -b_k(t) \times \langle S_k \rangle, \]

where \( \hbar \equiv 1 \) as we assume for the rest of the paper.

It is simple to show that momentum independent modulations of the quasiparticle energy are irrelevant. Indeed, consider such a time-dependent modulation in impulsive form,

\[ \varepsilon_k(t) = \varepsilon_k^0 + V \Delta t \delta(t), \]

where \( \varepsilon_k^0 \) is the equilibrium dispersion relation and \( V \Delta t \delta(t) \) is the strength of an impulsive potential that couples to the total charge. Integrating the equation of motion in a small interval of time \( \Delta t \) we find that after the impulse the pseudospins obey

\[ \langle S_k(t) \rangle = \langle S_k \rangle^0 - \hat{z} \times \langle S_k \rangle^0 V \Delta t, \]

where \( \langle S_k \rangle^0 \) is the equilibrium pseudospin before the impulse.
a consequence, uniform (momentum independent) drives as the one in Eq. (8) can be eliminated from the outset. For the same reason we can assume that \( \mu \) is time independent and corresponds to the equilibrium chemical potential before the drive. Also, for the same reason, long-range Coulomb interactions do not play any relevant role for this class of drives. Long-range Coulomb interactions may become relevant if one considers drives that couple to the density operator at finite momentum which is beyond the scope of this work.

### III. DRIVING MECHANISMS OF THE SUPERCONDUCTING CONDENSATE

In this section we discuss different mechanisms that can be used to drive the superconductor out of equilibrium. Contrary to previous works in the context of interaction quenches in ultracold-atomic systems [11–13,26,27], we shall consider a periodic time-dependent perturbation acting over a long time.

#### A. Density of states driving

The study of the coupling of phonons to the spontaneous amplitude mode of the order parameter started several decades ago in relation with 2H-NbSe₂ [16,17,28,29]. In this material a charge-density wave appears at an ordering temperature \( T_{\text{CDW}} = 33 \text{ K} \) above the superconducting critical temperature, \( T_{\text{c}} = 7 \text{ K} \). The CDW partially gaps the Fermi surface in such a way that a Raman phonon, which drives the CDW, strongly modulates the density of states (DOS) at the Fermi level available for superconductivity. As a consequence, at low temperatures in the superconducting phase the equilibrium amplitude of the order parameter is strongly dependent on the lattice coordinate, which we will denote as \( u \).

Since the relevant phonon is Raman active it can be launched impulsively [30,31]. Assuming that the light pulse is applied at \( t = 0 \) and that damping is negligible, the phonon coordinate obeys \( u(t) = u_0 \Theta(t) \sin(\omega_d t) \) where \( \Theta(t) \) is the Heaviside step function and \( \omega_d \) denotes the driving frequency (in this case the phonon frequency). A change in the DOS can be introduced as a change in the Fermi velocity which corresponds to a time dependence in Eq. (1) of the form \( \xi_k(t) = \xi_k^0[1 + \beta(t)] \) with \( \beta(t) \propto u(t) \). Because, as discussed above, dynamical terms that couple to the total density are irrelevant, we can add a time-dependent term proportional to the chemical potential so that the relevant perturbation is given by

\[
\xi_k(t) = \xi_k^0[1 + \beta(t)].
\]  

This corresponds to a change in the DOS, \( N(t) = N_0/[1 + \beta(t)] \), where \( N_0 \) is the equilibrium value. These equations show that by exciting the Raman phonon one can induce periodic oscillations in the DOS which will take the superconductor out of equilibrium.

At this point it is important to emphasize the following: the thermodynamic BCS gap equation for a mechanism with cutoff frequency \( \omega_{\text{D}} \),

\[
\Delta_0 = 2\omega_{\text{D}} e^{-1/N_0\lambda},
\]

suggests that a change in the DOS can be absorbed in a change in \( \lambda \). Such an assumption has been done in the past. However, in the present formalism, changing \( \lambda \) and changing the DOS through Eq. (10) are distinct drives—drive in the \( x \) component or in both the \( x \) and \( z \) components of \( b_k(t) \), respectively—that strictly speaking lead to distinct dynamics. Nevertheless, we find that at long time scales the results are qualitatively similar and the existence of the Rabi-Higgs mode and all physical consequences remain unchanged for both mechanisms.

Other alternative excitations that can lead to a time-dependent DOS, for instance by using impulsive stimulated Raman scattering (ISRS) and THz radiation pulse, are discussed in Appendix A.

#### B. Phonon-assisted coupling constant driving

An interesting type of driving consists of phonons that can modulate the coupling constant. We will refer to this case as \( \lambda \) driving. We propose that \( \lambda \) driving can be realized in unconventional superconductors, in particular Fe-based superconductors.

A compilation of several experiments in different materials of the Fe-based superconductors family [32–36] show that the critical temperature (and therefore the condensation energy) is very sensitive to the anion height from the Fe layer, in particular for FeSe (see Appendix B). Similar conclusions can be drawn from scanning tunneling microscopy experiments [37] and magnetic penetration depth measurements [38].

Theoretically [39–44] one finds in these materials that the magnetism is very sensitive to the anion’s height. This can be naturally explained if the system is close to a Stoner instability controlled by the latter. Furthermore, for the magnetic mechanism expected for these materials, the pairing interaction is controlled by the magnetic susceptibility [45,46]. Therefore, it is quite natural to attribute a large fraction (if not all) of the anion’s height sensitivity in \( T_{\text{c}} \) to a modulation of the pairing interaction via the magnetic susceptibility.

Another feature that makes Fe-based materials particularly suited for our propose is that in practically all materials there is an \( A_{1g} \) phonon mode which involves the anion’s height coordinate. This mode can be launched in real time in a pump-probe experiment by a stimulated Raman process [3,30,31]. Therefore, we consider that the coupling constant is a function of the anion’s height which can be manipulated with laser impulses. We have estimated that the changes in \( \lambda \) are smaller than 10% for FeSe (Appendix B). Below we will study modulations of \( \lambda \) up to 10% for illustrative proposes but our main results are robust and visible for much smaller values. Thus, in the following we will consider that

\[
\lambda(t) = \lambda_0 + \frac{d\lambda}{dt} u(t) = \lambda_0[1 + \Theta(t)\alpha \sin(\omega_d t)],
\]

and we will take the parameter \( \alpha \in [0, 0.1] \).

#### C. Drive in ultracold atoms

The manipulation of the Hamiltonian parameters in ultracold atoms is very well known [10] and therefore we mention it here very briefly. Both \( \lambda \) driving and DOS driving can be technically achieved in ultracold atoms.

The interaction between fermions can be controlled through a time-dependent magnetic field that modulates a Feshbach resonance [8] or through optical control which is
much faster [9]. Very recently, Ref. [7] has introduced a different way to modulate the order parameter which appears well suited for our propose. This method involves generating Rabi oscillations between one of the two fermionic states participating in pairing and a third state.

To modulate the DOS one can consider fermions moving in an optical lattice [10]. It is possible, then, to modulate the depth of the potential well in time as has already been done for bosonic systems [47–49]. Such drive modulates the DOS via the change in the hopping integral.

**IV. LINEAR RESPONSE**

As mentioned above, even though the DOS driving and λ driving are different, the main physical results are very similar. Thus, in the remainder of this work we analyze the λ driving in further detail.

As a warmup exercise we compute the linear response of the superconductor to a harmonic λ drive with adiabatic switching. In the following, we denote the corrections that are linear in the perturbation by a superscript “1,” that is \( b_k = b_k^0 + b_k^1 \exp[i(\omega - i\delta)t] \) for \( b_k^0 \) and \( b_k^0 + b_k^1 \exp[i(\omega - i\delta)t] \) with \( \alpha = \lambda_0 + \lambda_1 \exp[i(\omega - i\delta)t] \) where \( \delta \) is an infinitesimal positive quantity. The linearized equation of motion becomes

\[
i \omega S_k^0 = -b_k^0 \times S_k^0 - b_k^0 \times S_k^1,
\]

with solution

\[
S_k^1 = \frac{b_k^1 \xi_k}{2\sqrt{\lambda_0^2 + \xi_k^2}[(\omega - i\delta)^2 - 4(\lambda_0^2 + \xi_k^2)]} \left( \frac{-2\xi_k}{2\lambda_0} \right).
\]

Notice that since we are considering an s-wave superconductor there is no momentum dependence of the \( x \) component of pseudomagnetic field \( (b_k^1 \equiv b_x^1) \). Therefore, assuming a particle-hole symmetric DOS the only nonzero component after summing over \( k \) is

\[
S_i^{1x} = \chi_{\Delta,\Delta}(\omega)b_i^{1x},
\]

where the quantity on the left side was defined below Eq. (5) and we introduced the bare susceptibility,

\[
\chi_{\Delta,\Delta}^0(\omega) = -\sum_k \frac{\xi_k^2}{\sqrt{\lambda_0^2 + \xi_k^2}[(\omega - i\delta)^2 - 4(\lambda_0^2 + \xi_k^2)]}.
\]

The imaginary part of Eq. (16) is given by

\[
\text{Im}[\chi_{\Delta,\Delta}^0(\omega)] = -\text{sgn}(\omega) \frac{\pi \rho}{4} \sqrt{1 - \left( \frac{2\Delta_0}{\omega} \right)^2},
\]

for \( 2\Delta_0 < |\omega| < \omega_D \) and zero otherwise, while the real part can be obtained from the Kramers-Kroning relation

\[
\text{Re}[\chi_{\Delta,\Delta}^0(\omega)] = \frac{1}{}\Pi \int_{-\infty}^{\infty} \text{Im}[\chi_{\Delta,\Delta}^0(\omega')] d\omega'.
\]

Furthermore, we can write

\[
b_i^{1x}(t) = 2\Delta_1(t) = 2[\lambda_1(t)S_i^{0x} + \lambda_0 S_i^{1x}(t)],
\]

and by using Eq. (15) we obtain the pseudospin response for a positive frequency \( \omega \),

\[
S_i^{1x}(t) = \alpha \chi_{\Delta,\Delta}(\omega) S_i^{0x} e^{i\omega t},
\]

where we have defined

\[
\chi_{\Delta,\Delta}(\omega) \equiv \frac{2\lambda_0 \chi_{\Delta,\Delta}^0(\omega)}{1 - 2\lambda_0 \chi_{\Delta,\Delta}^0(\omega)}.
\]

Summing the response for \( \pm \omega \) and taking the real part, we get that the correction to the gap amplitude \( \Delta_1 \) induced by a \( \cos(\omega t) \) drive is

\[
\Delta_1 = \alpha [1 + \chi_{\Delta,\Delta}(\omega)] \Delta_0,
\]

and by using Eq. (15) we obtain that \( \Delta_0 \) is shown as a continuous line and dot dashed line respectively. The horizontal dashed line shows that the equation \( 1 - 2\lambda_0 \chi_{\Delta,\Delta}^0(\omega) = 0 \) is satisfied at \( \omega = 2\Delta_0 \). Bottom panel: the response determining the amplitude of the order-parameter oscillations excited by a \( \lambda \) drive. We use a cutoff frequency \( \omega_D = 2\Delta_0 \).

FIG. 1. Top panel: the real and imaginary part of susceptibility Eqs. (18) and (17) times \( \omega_0 \) is shown as a continuous line and dot dashed line respectively. The horizontal dashed line shows that the equation \( 1 - 2\lambda_0 \chi_{\Delta,\Delta}^0(\omega) = 0 \) is satisfied at \( \omega = 2\Delta_0 \). Bottom panel: the response determining the amplitude of the order-parameter oscillations excited by a \( \lambda \) drive. We use a cutoff frequency \( \omega_D = 2\Delta_0 \).
thermalization at infinite temperatures at very long times. However, as we will show below, going beyond linear response this does not occur in the present system, a fact that can be attributed to the integrability of the model \cite{51}. In an open system in contact with a thermal bath one expects that the energy will be transferred to the bath, usually the lattice phonons in a solid-state state superconductor. However, at low temperature these times can be very long \cite{1} leaving a large time window were the system effectively behaves as if were closed.

V. NONLINEAR RESPONSE AND RABI-HIGGS MODES

We now explore the behavior of the system when the drive acts for long times and/or the drive amplitude is not small.

A. Numerical results

We shall show the numerical calculation of $\Delta(t)$ beyond linear response. In our computations, for $t \leq 0$ the system is in equilibrium. The superconducting gap is set to $\Delta_0$ and assumed to be real. The zero-temperature pseudospin texture is given by Eq. (6). At $t > 0$ the drive switches on according to Eqs. (10) or (12) and the pseudospins evolve according to Eq. (7), which in turn will change the gap $\Delta(t)$ and the local fields $b_i(t)$. We take a set of $N$ pseudospins uniformly spaced in $\xi$ within a band of width $W = 40\Delta_0 = 2\omega_D$ which equals twice the cutoff frequency ($\omega_D$, for conventional superconductors). The fourth-order Runge-Kutta method was used to numerically integrate the spin equations of motion with $N = 4 \times 10^4$. If the initial order parameter is real (a gauge choice), then the symmetry of the problem dictates that it remains real at all times in both the case of $\lambda$ and DOS driving. Hence, we set Eq. (5) to be zero.

We assume a driving amplitude that is constant in time. In the case of phonon assisted driving, this may appear unrealistic as real phonons will have damping. More so in the presence of the superconductor where energy will be transferred from the phonon to the superconductor. However, in experiments this can be compensated by periodically applying pulses with a periodicity which is a multiple of the phonon period in order to restore the lattice oscillation to the original amplitude. Obviously, in the case of ultracold atoms, electronic ISRS (Appendix A 1) or THz (Appendix A 2) driving this problem does not arise.

In the following we present the results for the $\lambda$ driving of Eq. (12) in detail. We use parameters adequate for the phonon-assisted version in FeSe materials (Sec. III B). Qualitatively equivalent results were found with the DOS driving and with other parameters.

The dynamics of $\Delta(t)$ is shown in Fig. 2 for $\omega_d = 4\Delta_0$ and several values of $\alpha$. The first feature to notice is that despite that the drive can give energy to the superconductor indefinitely, the average order parameter decreases from the equilibrium value but it is not totally suppressed, i.e., the system is not driven to infinite temperature.

In Fig. 2 (and Fig. 4 below), the drive frequency corresponds to an oscillation that is too fast to be resolved on the scale of the figure and leads to the black filled regions. The system essentially synchronizes with the drive. Both features, synchronization and absence of heating, are expected \cite{51} for an integrable system as the present one. On top of that, the amplitude of the gap shows slow oscillations. This nonlinear low-frequency mode is our main result and will be dubbed the Rabi-Higgs mode.

The period of the Rabi-Higgs mode decreases with increasing $\alpha$ as will be discussed in detail below. It is important to realize that these driven modes are very different from the spontaneous Higgs modes reported before \cite{11–13} which have a different frequency and do not show a similar sensitivity to the drive.

In order to gain more information on the oscillations we perform a fast Fourier transform analysis. This is shown in the insets of Fig. 2. We find that for all studied $\alpha$ there are two fundamental frequencies in the dynamical response of $\Delta(t)$. The first is the drive frequency $\omega_d$ and the second one corresponds to the Rabi-Higgs frequency $\omega_R$ with period $\tau = 2\pi/\omega_R$. Moreover, satellite peaks can be observed at $\omega_d \pm n\omega_R$ with $n = 1, 2, 3$ are observed.

We now discuss the amplitude of the order-parameter oscillations at frequency $\omega_d$. Because of the Rabi-Higgs mode, the amplitude of the oscillations (around some average value) is not constant leading to the variable width of the black regions in Fig. 2. In order to roughly take into account this effect, we define an effective time-dependent amplitude in a long steady-state part of the dynamics denoted as $\Delta_1^{\min}$ and $\Delta_1^{\max}$, respectively.

In practice, they are the minimum and maximum widths of the black regions in Fig. 2 at long times. Figure 3 shows the two amplitudes as a function of the strength of the drive $\alpha$. The solid line is the linear-response result of Eq. (22) with $\omega = \omega_d$. We see that the latter predicts the correct magnitude of $\Delta_1$ and it is quite close to $\Delta_1^{\min}$, i.e., linear response.
fixes the scale of the amplitude of the oscillations at the drive frequency even far from its strict range of strict validity (small $\alpha$ and short times). Linear response cannot explain the difference between $\Delta_{\min}$ and $\Delta_{\max}$ which is an exquisitely nonlinear effect. Indeed such difference is associated with the appearance of a new frequency which is clearly an effect beyond linear response. This can be also seen from the inset of Fig. 3 where it is clear from the numerical results that for small $\alpha$ the difference between the two amplitudes is approximately quadratic in $\alpha$. There is a small mismatch of the slope at small $\alpha$ but this is not surprising as even in that regime we are not strictly in the conditions of validity of linear response since the switching was not adiabatic and the time of the measurement was not small. What is important for our purpose is that the overall scale is well predicted. As a further check that the scale of the amplitudes are determined by the linear-response’s susceptibility we change the drive frequency towards the Higgs resonance of Fig. 1. As expected, as $\omega_d \to 2\Delta_0$ the amplitude increases as is shown in Fig. 4.

In the remainder of the paper, we present simulations only for a drive frequency $\omega_d = 4\Delta_0$ but qualitatively similar results were found for $\omega_d = 6\Delta_0$ and $\omega_d = 8\Delta_0$.

B. NMR analogy

The drive induced Rabi-Higgs mode originates in a resonant phenomenon for pseudospins analogous to Rabi oscillations in NMR experiments for usual spins as we shall demonstrate below.

In the presence of a static magnetic field $B_0$ any spin or magnetic moment precesses with the Larmor frequency $\omega_L$, which is proportional to $B_0$. If a small alternating magnetic field of amplitude $B_1$ is applied with a frequency $\omega = \omega_d$ in a plane perpendicular to $B_0$, the spin experiences Rabi oscillations with a new low frequency which is proportional to $B_1$ [52],

$$\omega_R = \gamma B_1,$$

where $\gamma$ is the gyromagnetic ratio.

In the case of the pseudospins the role of the static magnetic field is played by the mean field $b_k$ and the Larmor frequency is $\omega_L = 2\sqrt{\xi_k^2 + \Delta_0^2}$. This is nothing but the energy needed to create two Bogoliubov quasiparticles as shown schematically in Fig. 5. There is a family of pseudospins, labeled $S_k$, for which the resonance condition

$$\omega_d = 2\sqrt{\xi_k^2 + \Delta_0^2}$$

is satisfied. We will refer to this family as resonant Anderson pseudospins (RAPs). The pseudomagnetic field acting on the RAP set is

$$b_k(t) = 2(\Delta(t), 0, \xi_k).$$

As a first approximation we schematize the time dependence of the field by

$$\Delta(t) = \tilde{\Delta} + \Delta_1 \cos(\omega_{rt}),$$

FIG. 5. Schematics of the proposed experiment. A drive at frequency $\omega_d$ enters into resonance with a set of Bogoliubov quasiparticles having energy $\pm E_k = \pm \sqrt{\xi_k^2 + \Delta_0^2}$. The RAP set is indicated by the red dots in this one-dimensional cut and corresponds to a line for the case of a two-dimensional superconductor.
FIG. 6. (a) Dependence of low-energy mode frequency $\omega_R$ on $\Delta_1$ both extracted from the numerical simulations. The line corresponds to the prediction of Eq. (28) assuming a constant amplitude $\Delta_1$ and making the approximation $\bar{\Delta} = \Delta_0$. The error bars indicate the range $\Delta_1^{\text{min}} < \Delta_1 < \Delta_1^{\text{max}}$ (see text). (b) $\omega_R$ as a function of $\alpha$. The empty dots represent $\omega_R$ extracted from simulations while the solid line is the estimated value after Eq. (28) where $\Delta_1$ depends on $\alpha$ according to Eq. (22). The final expression is displayed in the Conclusions as Eq. (29).

where $\bar{\Delta}$ is the average value of the order parameter in the steady state (similar but not necessarily equal to $\Delta_0$) and $\Delta_1$ is a constant oscillation amplitude at the frequency of the drive and whose magnitude can be approximated by linear response. We now decompose the time-dependent part of the pseudomagnetic field (which is directed along $x$) in a component parallel to the time independent pseudofield and another component perpendicular to it. The perpendicular component is

$$b_{k}^\perp(t) = 2\Delta_1 \cos(\omega_d t) \sqrt{1 - \left(\frac{2\bar{\Delta}}{\omega_d}\right)^2}.$$  

According to Eq. (23), this time-dependent magnetic field produces Rabi oscillations of the RAPs with frequency

$$\omega_R = \bar{\Delta} \sqrt{1 - \left(\frac{2\bar{\Delta}}{\omega_d}\right)^2},$$  

proportional to $\Delta_1$. To check that $\omega_R$ arises from this underlying mechanism we extracted, for each value of $\alpha$, $\Delta_1$, and $\bar{\Delta}$ from the full numerical solution for $\omega_k$ and compared it with Eq. (28). Figure 6(a) shows the plot of $\omega_R$ as a function of $\Delta_1$. The minimum ($\Delta_1^{\text{min}}$) and maximum ($\Delta_1^{\text{max}}$) of the order parameter are represented with a horizontal error bar. There is good agreement with the simulations with an effective fixed $\Delta_1$ within its physical range of variation. Figure 6(b) shows the frequency of the Rabi-Higgs mode as a function of the drive strength. The line is Eq. (28) where the value of $\Delta_1$ chosen is the one predicted by linear response, Eq. (22). This approximation appears to be surprisingly good as seen in the figure.

C. Physical consequences of Rabi-Higgs modes: Population inversion

We now analyze how the Rabi-Higgs mode affects different observables. In Fig. 7 the time dependence of the pseudospin texture is shown. Pseudospins are labeled by the quasiparticle energy and shown as a function of time with negative times representing the equilibrium situation. Colors encode the projection of the pseudospins in the different directions. Notice that $\langle S^z_k \rangle$ fluctuations, which determine the imaginary part of $\omega_k$, are odd so they cancel when summed over $k$. Therefore, only amplitude fluctuations enter into play in $b_k(t)$.

The drive frequency is visible in the dynamical response of all projections. Its period $T = 2\pi/\omega_d$ corresponds to the vertical features indicated by black arrows in the top panel of Fig. 7. We find that most of the pseudospins only precess, with small fluctuations, around their equilibrium values. In contrast, RAPs lying at $\xi_k$ and indicated by the white horizontal arrows respond strongly to the drive and visit the whole Bloch sphere in the Rabi period $\tau$, which is indicated by the white horizontal double arrow in the bottom panel of Fig. 7—it coincides with the period of the Rabi-Higgs mode. As in the analogous NMR experiment, RAPs perform slow oscillations from the equilibrium position to the antipode in the Bloch sphere and rapid rotations around the axis joining these points. Such direction corresponds approximately to the equilibrium field $\hat{b}_k = 2\Delta_0 \hat{x} + 2\xi_k \hat{z}$ and is a function of $\omega_d$ through the resonant condition Eq. (24).
increasing $\alpha$ (i.e., $n_\alpha$ fluctuation for these resonant states). The occupation values at equilibrium are indicated by continuous blue lines and denoted as $t < 0$.

Notice that the observed phenomena require that a macroscopic number of pseudospins synchronize due to interactions. Thus the RAP set is not limited to the only ones corresponding to a gapless regime. For large enough $\alpha$, the details of the gap dynamics become more complicated and will be presented elsewhere. On the other hand, for large enough $\alpha$, we find a far-from-equilibrium phase of gapless superconductivity.

Figure 9(a) shows, in logarithmic-linear scale, the gap dynamics for different values of $\alpha$. The order parameter first increases and then goes to zero exponentially in time showing oscillations with the drive frequency. As is shown in Fig. 9(b), inside this gapless phase, the $x$ component of pseudospin texture $\langle S_{x,k} \rangle$ is nonzero but its average is zero. It reflects the fact that there are pair correlations between fermions, for instance $\langle c_{x,k}^\dagger(t)c_{x,k}(t) \rangle \neq 0$ although $\Delta(t) = 0$, a consequence of a full pseudospin dephasing. In order to characterize this state, we define the quantity $\gamma = \beta \sum_k \langle S_{k}^x \rangle^2$ where we choose $\beta$ such that $\gamma = 1$ at $t = 0$ (equilibrium). While the superconducting gap (average of $\langle S_{k}^x \rangle$ essentially) goes to zero for long times, $\gamma$ remains as a finite constant value as is shown in Fig. 9(c).

We should emphasize that, in contrast to previous works [12,26,53], this gapless regime is obtained under the presence of drive. In this sense we obtain a dynamical phase of gapless superconductivity by increasing the drive strength.

VI. APPLICATION OF RABI-HIGGS OSCILLATIONS TO THE GENERATION OF DYNAMICAL PHASE TRANSITIONS

The spontaneous (drive-free) dynamical phases of a fermionic condensate were reviewed in the Introduction. In the initial theoretical proposal [11–13] these phases were reached after an interaction quench. Despite intense theoretical studies focused in the peculiarities of ultracold atoms [54,55] such phases have not been found experimentally so far. The situation is even worst in solid-state systems where it is difficult to imagine how large interaction quenches can be
achieved. In this section we describe an alternative route to the generation of spontaneous dynamical phases based on the use of the Rabi-Higgs mode that can be applied to both ultracold atoms and solid-state systems.

The idea is to apply the drive for an amount of time $t^* < \tau$, i.e., smaller than a full Rabi cycle, in such a way that only a fraction of the population inversion of RAPs has been achieved. Then the system is allowed to evolve spontaneously without the drive. In the case of phonon-assisted drives generated with ISRS the lattice motion can be stopped applying a second pump pulse delayed from the first one by a time equal to a half integer times the drive periodicity, i.e., $(n + \frac{1}{2})\tau$, as can be easily seen solving the equations of motion for a driven harmonic oscillator. On the other hand, THz drives and drives in ultracold atoms can be turned on and off at will with standard experimental setups.

With the above scheme the Rabi-Higgs mode dynamics can be seen as the preparation of an initial condition to the subsequent free (unperturbed) BCS time evolution. Different protocols have been used in previous works to set up the initial condition and only oscillations with damping and overdamped dynamics of $\Delta(t)$ were obtained \cite{20,53}. In contrast, here we shall demonstrate the possibility to access all dynamical phases of the order parameter. In our proposal we can induce a full population inversion in the pseudospin texture as an initial condition and it is precisely this that makes the persistent oscillation of the superconducting gap also accessible.

**Numerical simulation**

We test this idea numerically with $\lambda$ drives switching on the drive at $t = 0$ as before, and switching it off at $t = t^*$ during the first period of the Rabi cycle following the subsequent spontaneous out-of-equilibrium dynamics. Remarkably, for moderate values of $\alpha$ we obtain basically two dynamical phases, corresponding to relaxation or persistent oscillations of the order parameter, depending on the fraction of the Rabi cycle that has been completed. This is parametrized by $t^*/\tau$.

Figure 10 shows the gap evolution for two values of the drive, $t^*/\tau$. For small $t^*/\tau$ the gap asymptotically approaches a constant value $\Delta_\infty < \Delta_0$ exhibiting the well-known oscillatory behavior with $1/\sqrt{T}$ decay and frequency $2\Delta_\infty$ \cite{11–13}. For intermediate values of $t^*/\tau \sim 0.5$ the gaps show spontaneous persistent oscillations between the two extrema $\Delta_\pm$ as reported before for interaction quenches \cite{11–13}. The dynamical phase diagram for the proposed protocol with the long-time gap values of the gap and $\Delta_\pm$ values is summarized in the bottom panel of Fig. 10. In the gray region spontaneous persistent oscillations are observed while in the white region one sees the decaying oscillations.

In the previous section, for large values of $\alpha$, we have obtained a gapless phase under the presence of the drive [see Fig. 9(a)]. It is natural to ask what is the fate of this state if the drive is turned off at a time $t^*$ in the gapless regime. We checked that if $\Delta(t^*)$ is small enough the gapless regime remains. A similar result has been obtained before in a study of the response of a superconductor to a short but intense pump pulse in which $\Delta(t)$ is strongly suppressed before the drive-free evolution \cite{53}.

It is interesting to point out that the dynamical phase diagram of Fig. 10 is consistent with the Lax reduction method used in the literature \cite{13,56} (see Appendix C).

**VII. CONCLUSIONS**

We have shown that a superconductor subject to a periodic harmonic drive in the collisionless regime performs Rabi-Higgs oscillations with a frequency that can be controlled by the strength of the drive. We argued that there are several solid-state and ultracold atom routes to realize the drive.

Periodic drives have been considered before \cite{14,15,57}. Our results go beyond previous computations of third-order susceptibilities \cite{14,18,57} by considering the full dynamical nonlinear response beyond perturbation theory. Reference \cite{15} finds a slow oscillation of the superconducting order parameter under a drive, in a model in which superconductivity is close to a charge-density-wave instability. Then the system oscillates periodically between different orders with a frequency much slower than the drive. Our results apply instead to the more general case of a BCS superconductor which is not close to another order and the oscillations are between an equilibrium superconductor and a highly excited superconductor.

In real solid-state and cold-atom systems, terms in the Hamiltonian beyond BCS mean field will produce collisions and tend to relax the out-of-equilibrium populations. It is hard to predict in advance which system can fulfill the requirement that the coherence time is long enough to see the Rabi-Higgs oscillations. In general low pump fluence will reduce heating but produce slower oscillations which may require longer
coherence times to be seen. Therefore, a tradeoff should be found which will depend on the specific pump mechanism and material. In the present mechanism the drive frequency has to be larger than 2\(\Delta_0\) but be of the same order. Keeping \(\alpha\) and the ratio \(\omega_d/\Delta_0\) fixed, the Rabi frequency scales linearly with the gap [cf. Eqs. (22) and (28)]:

\[
\omega_R = \alpha |1 + \chi_{\Delta,\Delta}(\omega)| \Delta_0 \left(1 - \left(\frac{2\Delta_0}{\omega_d}\right)^2\right).
\]

thus lowering the gap makes the Rabi oscillation longer, \(\tau \sim 1/\Delta_0\). On the other hand, coherence times usually scale faster with quasiparticle energy. For example in a Fermi liquid, decay due to electron-electron interaction scales as \(1/\omega^2\). Therefore, working with materials with small gaps can be beneficial as coherence times can be made longer than the Rabi time one would like to measure. Taking into account that the system is not a Fermi liquid but a superconductor the situation is, of course, even better. In this regard it is encouraging to notice that Rabi oscillations are not unprecedented in the solid state but have a long history in semiconductors [58–60] and are part of the modern toolbox of solid-state quantum technologies [61–64].

We have discussed drives which preserve the symmetry of the lattice and couple to the Higgs modes. It is also possible to apply drives that do not preserve the symmetry of the lattice and couple to the Higgs modes. It is also possible to apply drives that do not preserve the symmetry of the lattice and couple, for example, to \(B_{1x}\) and \(B_{2x}\) charge fluctuations in the case of a square lattice. These drives will lead to Rabi-charge oscillations entirely analogous to the Rabi-Higgs oscillations. An advantage of these modes is that being Raman-like they modulate the optical properties of the material as the analogous spontaneous modes do [2,3]. Therefore, they can be studied by optical means for example by pumping in the THz range and probing through differential reflectivity with visible light.

As an application we have shown how the induced far-from-equilibrium state can be used to observe the so far elusive spontaneous dynamical phases predicted by theory [11–13]. The original proposals were based on a very strong perturbation applied in a short time. Instead, our proposal is based on a gentle perturbation applied for a moderate time exploiting the properties of collisionless or nearly collisionless condensates. This opens a window of opportunity to make these phases accessible even in a solid-state setting.

Our work is an example of quantum control of a condensate wave function. The population of a set of quasiparticles is less condensates. This opens a window of opportunity to make

\[
\frac{\omega_d}{\Delta_0} \lesssim 0.1.
\]

APPENDIX A: ELECTRONIC DRIVING MECHANISMS

1. ISRS driving

Electronic Raman-active excitations of the condensate [2,3] can be used to drive a superconductor through an ISRS process analogous to ISRS for phonons [30] or magnons [67]. Neglecting absorption at the pump-laser frequency, the total Hamiltonian for this process is \(H = H_0 + H_R\) with \(H_0\) the equilibrium BCS Hamiltonian and \(H_R\) the electronic Raman Hamiltonian [2,3], which is given by

\[
H_R = -2 \sum_{X,k} v_X(t) f_k^X S_k^x,
\]

\[
v_X(t) = -\frac{1}{2} \frac{\partial}{\partial N_X} \langle N_X \rangle \frac{\partial}{\partial \omega_X} \langle \omega_X \rangle.
\]

Here \(E(t)\) is the light electric field of the pump laser, \(v_X = \chi_{\omega_X} / \partial N_X\) is the Raman tensor for electronic Raman scattering in symmetry \(X\), and \(\omega_L\) is the laser carrier frequency. The description in terms of the electric field (instead of the vector potential as in next subsection) emphasizes the relation with the optical properties at the energy of the pump, \(\omega_L\).

Restricting to a tetragonal layered material, the most relevant symmetry functions are

\[
\begin{align*}
 f_k^{A_{1g}} &= \frac{1}{2} [\cos(k) + \cos(k)], \\
 f_k^{B_{1g}} &= \frac{1}{2} [\cos(k) - \cos(k)], \\
 f_k^{B_{2g}} &= \cos(k) \sin(k).
\end{align*}
\]

Notice that we have not included \(f_k^{A_{1u}} = 1\) as it leads to driving by the total number operator which is irrelevant for the reason explained in Sec. II. Equation (A1) shows that in this case the laser electric field acts as a time-dependent potential acting on charge excitations with different symmetries. Usually pulses of about 50 fs can be produced which modulate an IR or visible laser. Mansart et al. have shown [2] that one such pulse induces a fluctuation of the condensate at a frequency close to \(2\Delta\) by an ISRS process. Also here it is possible to adjust \(\omega_L\) to a window of low absorption to minimize heating.

The charge fluctuation after one pulse decays very rapidly because of dephasing of excitations with different frequency. However, as mentioned in [3], one can excite the material with a periodic sequence of pulses so that the excitations that match the periodicity of the pump are reinforced and other excitations are suppressed. It is precisely the dynamics of this kind of excitation that is described in detail in Sec. V. One advantage of this method is that it is very easy to detect the response of the condensate through the modification of the optical properties at optical or other frequencies [2,3].

Selection rules for ISRS are similar to the one of spontaneous Raman scattering except that in the former \(v_X\) involves the same polarization of the electric field at both sides of
the Raman tensor in Eq. (A1) while in the latter two different fields appear related to incoming and outgoing photons. Therefore, in the notation of Table I of Ref. [68], polarizations $xx, yy,$ and $x'x'$ are accessible in ISRS while $xy$ is not. These “parallel” polarizations excite the $A_{1g}$ symmetry modes plus other modes. By symmetry, and using the same arguments as before, the $A_{1g}$ part of the drive can be taken to the DOS-driving form of Eq. (10) plus subleading terms with more complicated structure along the Fermi surface. This is obvious in the case of a lattice model with only nearest-neighbor hopping since the $A_{1g}$ symmetry function is proportional to the dispersion relation. For reasons of simplicity, in our simulations we concentrate on modes that preserve the symmetry of the lattice but we expect that for modes of lower symmetry very similar physics arises.

2. Direct THz drive

Matsunaga et al. have shown [4,5] that THz radiation pulses can produce oscillations of a superconducting condensate. They interpreted this result as due to the coupling of the Higgs amplitude mode to the THz electric field. However, this interpretation has been disputed by Cea, Castellani, and Benfatto [6] who argued that the response is dominated by charge fluctuations similar to the ones of the transient Raman experiment of Mansart et al. [2]. Notwithstanding, this kind of drive is very interesting since it is possible to apply a THz radiation with a frequency smaller than the gap in such a way that, to leading order, there is no direct excitation of quasiparticles and one expects much less heating than with an ISRS drive.

We introduce the coupling with the electromagnetic field through the Peierls substitution for carriers with charge $q \approx -|e|$ for electrons) in the underlying lattice model,

$$c_{i+r}^\dagger c_i \rightarrow c_{i+r}^\dagger c_i e^{-i\mathbf{A} \cdot \mathbf{r}},$$

(A3)

where $\mathbf{A} \equiv q \mathbf{A}/(\hbar c)$ and we assume that the vector potential $\mathbf{A}$ is uniform on the scale of the sample giving rise to the electric field, $E = -\mathbf{A}/c$.

Expanding up to second order in the vector potential we have that the Hamiltonian becomes $H = H_0 + H_T$ with

$$H_T = -2 \sum_k \left( \frac{\partial^2 \bar{\varepsilon}}{\partial k} \cdot \mathbf{a}(t) + \frac{1}{2} \frac{\partial^2 \bar{\varepsilon}}{\partial k_\mu \partial k_\nu} \alpha_{\mu\nu}(t) \alpha_{\nu}(t) \right) \mathbf{S}_k,$$

(A4)

where the sum over repeated indexes is implicit. The first term inside the parentheses represents the coupling of the vector potential to the paramagnetic part of the current. Consider $\mathbf{A}(t) = \hat{A} \Theta(t) \sin(\omega_T t)$. For $\omega_T < 2\Delta$ and not very strong disorder, there is no absorption from the paramagnetic part (the real part of the optical conductivity is zero) and there is no transfer of energy between the drive and the system. Since we are interested in processes where the drive and the system exchange energy, we consider only the second term. This behaves as a drive with a time dependence $\sim \cos(2\omega_T t)$. If we now consider a system with the same symmetries of the previous subsection and the $x'x'$ polarization so that $\hat{A}_x = \hat{A}_y$ we reach again a momentum dependence with the same symmetry as the dispersion relation (plus higher-order corrections). In analogy with the ISRS case, this can be taken to the DOS-driving form of Eq. (10). Notice that due to the similarity to the ISRS case it is quite tempting to excite in the THz to avoid heating but to use a probe at optical energies where it is easy to achieve high resolution in time.

**APPENDIX B: 1-DIVING MECHANISMS (FeSe case)**

For FeSe [34], the critical temperature changes from $T_c = 12$ K for $z = 1.457$ Å, to $T_c = 34$ K for $z = 1.427$ Å, where $z$ is the Se height from the Fe layer. This has been shown in experiments under pressure and it has being argued that the enhancement of $T_c$ should originate on the crystal structure because the total carrier density of the FeSe layer does not change with pressure. Thus, this material is particularly appealing to modulate the order parameter with a time-dependent lattice distortion. As an order of magnitude estimate, the numbers quoted above amount to a rate of change of $T_c$ with the anion height of 7 K/Å.

Exciting FeSe with a 1.5-eV infrared light pulse and a fluence of about 0.46 mJ/cm$^2$, Gerber et al. found that an amplitude of ±0.25 Å is achieved for the $A_{1g}$ phonon that oscillates at a frequency $\omega_0 = 5.3$ THz [69]. In the static limit, such a displacement would correspond to a variation of $T_c \sim \pm 2$ K which translates in a variation of about 10% in the zero-temperature equilibrium gap. Since the changes in the DOS within band theory are not of that order, it is natural to assume that such dramatic variations are due to phonon induced changes in the pairing interaction as explained above. To estimate the change in $\lambda$ one can use the BCS thermodynamic gap function to obtain $\delta \lambda/\lambda = N_0 \delta \Delta/\Delta$ where $N_0$ is the density of states at the Fermi level. Since $N_0 \lambda < 1$, the changes in $\lambda$ are smaller than 10%.

The quoted fluence is close to the one that was used in Ref. [2] (0.3 mJ/cm$^2$) to observe oscillations of the superconducting condensate in cuprates ($T_c = 40$ K). Therefore, it is also reasonable to assume that a similar fluence in a FeSe superconductor will not destroy the superconducting condensate. Notice that the ISRS mechanism to launch the phonon does not require absorption. Therefore, the energy deposited in the sample could be further minimized by tuning the laser excitation energy to a transparent energy region of the material, allowing us in principle to increase the amplitude of the oscillation without heating.

In FeSe there is a dramatic enhancement of the equilibrium superconducting $T_c$, as sample thickness is reduced. $T_c$ changes from 8 K in bulk to 77 K in a monolayer grown on SrTiO$_3$ which corresponds to a similar variation of the gap parameter. Thus, by controlling the film thickness and assuming the phonon frequency does not change much, we have a wide range of the ratio $\omega_0/2\Delta_0 \sim 1$–10 which is experimentally accessible and worthwhile to explore in numerical simulations.

**APPENDIX C: LAX REDUCTION ANALYSIS**

We verified the dynamical phase diagram of Fig. 10 by using the Lax reduction method [13,56]. The integrability of the model in the absence of drive implies that the frequency spectrum that determines the spontaneous evolution of the
order parameter is determined by a set of integrals of motion. The latter can be evaluated at any time with the so-called Lax vector \[12,13\] defined as a function of an auxiliary parameter \(u\),

\[
L(u) = z + \lambda \sum_k S_k \frac{u - \xi_k}{u^* - \xi_k^*}.
\]

The square of the Lax vector is a conserved quantity under time evolution with the unperturbed BCS Hamiltonian and therefore its roots (in the following Lax roots) are also conserved. Since the square of the Lax vector is non-negative, all roots should come in complex-conjugated pairs. Also, since 

\[
S_k^* = S_k^* - S_k^* - S_k^* - S_k^* \text{ hold at all times, where } k \text{ and } -k \text{ label the quasiparticle states with energy } \xi_k \text{ and } -\xi_k, 
\]

respectively, it is easy to see that if \(u\) is a Lax root \(-u\) is also a root (see top panels of Fig. 11).

In Ref. \[13\] it was shown that the dynamics of the order parameter is related to the number \(m\) of isolated pairs of complex-conjugated Lax roots. Indeed, \(\Delta(t)\) shows persistent oscillations at long times with \(k\) different frequencies if \(m > 1\) while \(\Delta(t) \to \Delta_\infty \) (damped oscillations) if \(m = 1\). Here, \(k\) is equal to the integer part of \(m/2\).

To check the numerical results we constructed the Lax vector and compute its roots from the instantaneous pseudospin texture at time \(t^*\) (see main text), i.e., with the initial condition of the subsequent free evolution. As an example, we show in the top panels of Fig. 11 the Lax roots in the complex plane for the same values of \(t^*/\tau\) that were used in top panels of Fig. 10. Consistently with the numerical results, for \(t^* = 0.1\tau\) we obtain one pair of isolated Lax roots \((m = 1, \text{damped dynamics})\) and for \(t^* = 0.5\tau\) we obtain \(m = 3\) (persistent oscillations).

For any value of \(t^*\) we have one pair of purely imaginary isolated Lax roots \(\pm u_{0}\) plus a continuum of doubly degenerated roots on the real axis. In the damped regime, the absolute value of the purely imaginary roots determine the asymptotic value of the superconducting gap, \(\Delta_\infty = u_0\) \[12\]. In the persistent oscillatory phase two extra pairs of isolated roots \(\pm u_r(\pm u_r^*)\) appear [see Fig. 11(b)]. We find numerically that \(\text{Re}(u_r)\) is close to \(\xi_k\) and the imaginary part is finite. Varying \(t^*/\tau\), the change of behavior of the damped regime is determined by \(\text{Im}(u_r) \to 0\). In the bottom panel of Fig. 11 we plot the imaginary part of these Lax roots \(u_r\) as a function of \(t^*\) which reproduces the phase diagram obtained numerically (cf. Fig. 10). We also show \(u_0\) vs \(t^*\) which accounts for the values of \(\Delta_\infty\) shown in Fig. 10.

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