

Superconductor-insulator transition driven by local dephasing

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We consider a system where localized bound electron pairs form an array of "Andreev"-like scattering centers and are coupled to a fermionic subsystem of uncorrelated electrons. By means of a path-integral approach, which describes the bound electron pairs within a coherent pseudospin representation, we derive and analyze the effective action for the collective phase modes which arise from the coupling between the two subsystems once the fermionic degrees of freedom are integrated out. This effective action has features of a quantum phase model in the presence of a Berry phase term and exhibits a coupling to a field which describes at the same time the fluctuations of density of the bound pairs and those of the amplitude of the fermion pairs. Due to the competition between the local and the hopping induced non-local phase dynamics it is possible, by tuning the exchange coupling or the density of the bound pairs, to trigger a transition from a phase ordered superconducting to a phase disordered insulating state. We discuss the different mechanisms which control this occurrence and the eventual destruction of phase coherence both in the weak and strong coupling limit.

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I. INTRODUCTION

The problem of interacting Cooper pairs and/or bosons, together with the possibility of a quantum control of long range phase coherence in low dimensional systems has received considerable interest since the experiments¹ on homogeneous lead and bismuth films, which exhibit a transition from a superconducting to an insulating phase as a function of thickness.

It is known, since the work of Abrahams et al.², that no true metallic behavior can be expected to be observed for 2D non-interacting electrons at $T=0$, because all the states will be localized by an arbitrarily small amount of disorder. When one includes Coulomb interaction, the situation is less clear, but the common belief is that a metallic phase should still not appear at $T=0$ in the presence of disorder - though no rigorous proof is available³. Yet, in the presence of attractive interactions, one expects a superconducting state both at $T=0$ and finite temperature, even in presence of a finite amount of disorder, due to its relative ineffectiveness for a transition of Kosterlitz-Thouless type (the degree of relevance being given by the Harris criterion⁴).

The existence of a superconducting state at $T=0$ in 2D systems is then considered to be directly linked to that of an insulating state, with no intermediate metallic phase present. One hence should be able to observe a direct superconductor-insulator transition (SIT) at zero temperature in 2D as a function of disorder, interaction strength, magnetic field or any other external parameter which can drive the system away from the superconducting phase.

There are different theoretical scenarios which are commonly discussed in connection with the salient features of such quantum phase transitions: i) dissipative mod-

els, considering a network of Josephson coupled superconducting grains shunted by resistors^{5,6,7,8,9,10}, where the competition between the Josephson intergrain coupling and the charging energy yields an increase of the phase fluctuations of the superconducting order parameter and hence leads to a phase disordered state. Dissipation thereby plays the role of suppressing quantum phase fluctuations and thus competes with the charging mechanism. ii) Bose-Hubbard models, where the superconducting phase is due to the Bose-Einstein condensation of charge-2e bosons and the insulating phase due to a proliferation of vortices and localization of pairs^{11,12,13,14}.

While the SIT was long thought to be a paradigm of the above theories, recent experiments have thrown some doubt about that. They revealed, what seems to be a low-temperature metallic state which is intertwined in such a SIT and thus requires a more appropriate theoretical description. In magnetic field tuned experiments on Mo-Ge samples¹⁵, in granular superconductors (Ga films¹⁶ and Pb films¹⁷), and in Josephson junction arrays¹⁸ a metallic phase has been observed. Moreover, the recent experiments of Kapitulnik et al.¹⁹, in which a metallic phase has been observed sandwiched between the superconductor and the insulating phase, suggest that perhaps two phase transitions accompany the loss of phase coherence in 2D superconductors : (i) a superconductor to a "Bose metal" and (ii) a "Bose metal" to an insulator. A "Bose metal" in this context, is thought of just a gapless non-superfluid liquid with metallic like transport²⁰.

These experimental results have led to reconsider the whole issue of the SIT. One recent proposal to handle such a new viewpoint of the SIT has been to reexamine the standard on-site charging model. By including nearest neighbor charging terms it was shown that

the resulting uniform Bose metal state lacks any trace of either phase or charge order²¹, due to a competition between the order parameters which describe the onset of charge order and that of phase coherence. A different point of view²² is, that in the quantum disordered regime a cancellation happens between the exponentially long quasi-particle scattering time and the exponentially small quasi-particle population, which ultimately leads to a finite dc conductivity. Finally, intrinsic as well as extrinsic sources of dissipation have been suggested as potentially relevant for the occurrence of a non-superfluid metallic phase in proximity to the superconducting phase^{19,23,24}. In this context, the dissipation is a relevant perturbation²⁵, which, depending on the strength of the coupling to the dissipative source, can drive the system from a *superconductor-to-insulator transition* to a *superconductor-to-metal transition*^{19,26}. Several other ideas supported a scenario where the system could break into superconducting and insulating "islands" weakly linked via percolating paths before going into a metallic phase, dominated by vortex dissipation²⁶. This "puddle" scenario matches with the view to describe the 2D superconductor-metal transition via superconducting islands embedded in thin metal films²⁷. Detailed analysis of the puddle-like model considered strongly fluctuating superconducting grains embedded in a metallic matrix²⁸, predicting a *metal-to-superconductor transition* with a metallic phase just above the transition dominated by Andreev reflections between the superconducting grains. Finally, recent investigations directed the attention to new phases with Bose metallic features, where the dissipation is dynamically self-generated, as in the quantum phase glass model, where disorder in the distribution of the tunnelling amplitudes and quantum fluctuations destroy phase coherence²⁰.

On similar lines as those dealt with in a great variety of such different approaches discussed in the literature, our aim here is to investigate a system where the breakdown of superconductivity situates itself in between the case of a "bosonic" and a "fermionic" mechanism for superconductivity suppression. We focus on determining the possible ground states for interacting Cooper pairs, in close relation to the classical notions of superfluidity and localization of bosons and, on a more general basis, how the phase coherence can be tuned to a phase disordered state whose transport properties may be unconventional. In particular, we consider within the framework of a boson-fermion model (BFM), a system where localized bound electron pairs (hard-core bosons) form an array of "Andreev"-like scattering centers coupled to a fermionic subsystem of itinerant uncorrelated electrons (quantum pair-exchange). This scenario goes beyond that of pure phase models widely discussed in the literature, since here one is dealing with bosonic degrees of freedom (localized bound pairs) as well as fermionic ones (itinerant electrons) for which, due to the emergence of pair correlations, one is dealing with both amplitude as well as phase modes. The possibility to tune from short

to long-range phase coherence arises from the following competing effects:

(i) on the one hand, the short range interaction between bosons and fermion pairs (holes) induces a local phase locking in a configuration with a quantum superposition between bosons and electron pairs, leaving the common phase undetermined.

(ii) on the other hand, the itinerancy of the electrons tends to lock and rigidly extend these initially arbitrary local phases. As a result, phase coherence develops over longer distances by suppressing the quantum fluctuations of the local phase, which will involve the dynamics of amplitude fluctuations.

Out of this competition one can recover either a superfluid state in the regime of a small scattering rate or a phase disordered state in the limit where the pair exchange dominates and the local quantum phase fluctuations do not allow for long-range phase coherence. As we shall show, the physics described by this scenario strongly depends on the concentration of fermions and bosons, on the coupling strength, as well as on temperature and magnetic fields. The purpose of this study is to analyze how a transition from a superconductor to an insulator may occur and discuss the features which can give rise to unconventional dynamics and, eventually, unconventional transport properties in the proximity of such a transition.

The outline of the paper is as follows. In Sect. II we will introduce the BFM and describe its main qualitative features and the phenomena which can be described by it. In Sect. III, we shall develop a path integral representation of this model and derive an effective coarse grained action, which, after integrating out the fermionic degrees of freedom, is able to describe the low energy dynamics of the phase and amplitude modes. In Sect. IV we will discuss the phase diagram based on the simplest approximation of such a coarse grained effective action in terms of a phase-only action and explore the transition from a superconductor to a phase disordered state as a function of the coupling and the density of bosons. In Sect. V, we will study the physical features which arise from the intrinsic Berry phase term present in our effective action and which arises from the hard core nature of the bosons, represented by quantum pseudo spin- $\frac{1}{2}$ variables. In section VI we compare the salient features of this BFM scenario with similar scenarios, such as the negative-U Hubbard model, the Bose-Hubbard model and Josephson junction arrays. In the Discussion, section VII, we review the main results obtained in this paper and indicate further developments planned for the near future.

II. THE BOSON-FERMION MODEL

The boson-fermion model (BFM) in recent years has attracted considerable attention as a model capable of capturing basic physical properties in many body systems with strong interaction, giving rise to the formation

of resonant pair states of bosonic nature inside a reservoir of fermions. Such a scenario was initially proposed by one of us (JR) as an alternative to the scenario of the hypothetical and yet to be experimentally verified *bipolaronic superconductivity*. It was meant to describe the intermediary coupling regime between the adiabatic and anti-adiabatic limits in polaronic systems, where an exchange between localized bipolarons and pairs of uncorrelated electrons can be assumed to take place (for a recent intuitive justification of such a scenario see for instance ref.²⁹).

As it has turned out, this boson-fermion scenario has a much wider range of applicability than that for which it was initially proposed and seems to apply to very different physical situation such as: hole pairing in semiconductors³⁰, isospin singlet pairing in nuclear matter³¹, d-wave hole and antiferromagnetic triplet pairing in the positive-U Hubbard model³² (and possibly also in the t-J model), entangled atoms in squeezed states in molecular Bose Einstein condensation in traps³³ and superfluidity in ultracold fermi gases induced by a Feshbach resonance³⁴.

The BFM is reminiscent of an anisotropic Kondo lattice model in terms of a pseudospin- $\frac{1}{2}$, but characterizing localized electron pairs instead of localized impurity spins as in the Kondo analogue. Its Hamiltonian is given by

$$H = (D - \mu) \sum_{i\sigma} c_{i\sigma}^{\dagger} c_{i\sigma} + (\Delta_B - 2\mu) \sum_i (\rho_i^z + \frac{1}{2}) \\ + \sum_{i \neq j, \sigma} t_{ij} (c_{i\sigma}^{\dagger} c_{j\sigma} + H.c.) + g \sum_i (\rho_i^+ \tau_i^- + \rho_i^- \tau_i^+)$$

The pseudo-spin operators $[\rho_i^+, \rho_i^-, \rho_i^z]$ denote the local bound electron pairs (bosons) and $[\tau_i^+ = c_{i\uparrow}^{\dagger} c_{i\downarrow}^{\dagger}, \tau_i^- = c_{i\downarrow} c_{i\uparrow}, \tau_i^z = 1 - \tau_i^+ \tau_i^-]$ the itinerant pairs of uncorrelated electrons. $[c_{i\sigma}^{\dagger}, c_{i\sigma}]$ stand for the creation and annihilation operators of the itinerant electrons (fermions) and g is the strength of the *boson* \Leftrightarrow *fermion pair* exchange interaction. The hopping integral for the itinerant electrons, which is assumed to be different from zero only for nearest neighbor sites, is given by t with a band half width equal to $D = zt$, z denoting the coordination number of the underlying lattice. The energy level of the bound electron pairs is denoted by Δ_B . The number of the ensemble of bosons and fermions being conserved, $n_{tot} = n_{F\uparrow} + n_{F\downarrow} + 2n_B$, implies a common chemical potential μ for both subsystems. $n_B, n_{F\uparrow, \downarrow}$ indicate the occupation number per site of the hard core-bosons and of the electrons with up and down spin states.

The exchange coupling between the bosons and the fermion pairs can be considered as an effective Andreev-like scattering leading to local states which are quantum superpositions of the form

$$|\psi_{loc}\rangle_i = \int d\phi_i [\cos(\theta_i/2) \cos(\phi_i) \rho_i^+ + \sin(\phi_i)] \times \\ [\cos(\phi_i) + \tau_i^+ \sin(\theta_i/2) \sin(\phi_i)] |0\rangle. \quad (1)$$

Such states evolve gradually out of the system of localized dephased bosons and essentially uncorrelated free fermions, which characterize the high temperature phase of this model, when the temperature is decreased below a certain $T^* \simeq g$ where resonant pairing (not bound pairs!) starts to be induced in the fermionic subsystem. These pair states have already features built in which are reminiscent of those which characterize Cooper pairing of fermions as well as superfluidity of bosons. The phases of the two coherent states, corresponding to the two subsystems, are the same and hence locked together, but are averaged over all angles as a consequence of the conserved particle number on any given site, $n_{tot} = 2$. Roughly speaking, the ground state of the system is then given by a product state $\prod_i |\psi_{loc}\rangle_i$ with $\cos\theta_i = 1$ and which exhibits no phase correlations on any finite length scale. Let us next consider the effect of fermion hopping between adjacent sites. This will give rise to density fluctuations on each of those individual sites and thus help to stabilize an arbitrary but finite average value of the phases $\{\phi_i\}$ over a finite length and time scale. In this way the localized bosons and fermion pairs acquire itinerancy^{35,36,37} which eventually leads to a superfluid state in both subsystems³⁸, provided that the effect of the local correlations between the bosons and the fermion pairs can be sufficiently diminished, but remaining still sufficiently strong to guarantee the formation of pairing in the fermionic subsystem. Achieving or not this situation will depend on the relative importance of the local exchange coupling versus the fermion hopping rate (given by the ratio g/t) and as well as on the concentration of the bosons, as shown by exact diagonalization study³⁹ on finite clusters of this BFM.

What we shall attempt in the present study is to describe this physics in terms of an effective action for the phase and amplitude fields of the bosonic fields. In order to achieve this we shall put the discussion on a level which is more familiar, namely that one of Josephson junction arrays and Bose-Hubbard models. For that purpose let us briefly sketch the analogy which exists (up to a certain point) between the BFM and those systems which have been widely discussed in the literature. A physically possible realization of this BFM scenario can be imagined in form of a network of superconducting grains embedded in a metallic environment and where the only mechanism of interaction between the grains and the fermionic background is that of Andreev reflections. Via such a mechanism an electron (hole) is reflected on the grain as a hole (electron) leaving behind a surplus of two holes (electrons) in the fermionic subsystem and of two electrons (holes) in the grain. If the grains are such that they have a large charging energy, the fluctuations of the number of pairs on them are energetically unfavorable and hence are largely suppressed. We then have a situation where the state of the grains switches essentially between zero to double occupancy with respect to the average occupation, any time an electron (hole) is reflected at the interface of the grain. Thus, the quantum dynam-

ics of the single grain can be directly represented by a pseudospin- $\frac{1}{2}$, in order to account for the doublets which represent the two possible states of the grains.

For such a possible experimental setup, the effective sites in the BFM have to be considered as defining a regular array of grains and having the same periodicity as the underlying lattice on which the fermions move with a hopping amplitude t . Moreover the size of the grains has to be such that it is much smaller than the distance between them. A pictorial view of such an experimental set up is given in Fig. 1. The analogy between the BFM and the array of superconducting grains which scatter pairs of fermions in a metallic matrix via Andreev like reflections, may ultimately serve as an experimental device on which to test and analyze the theoretical issues which will be discussed in this paper.

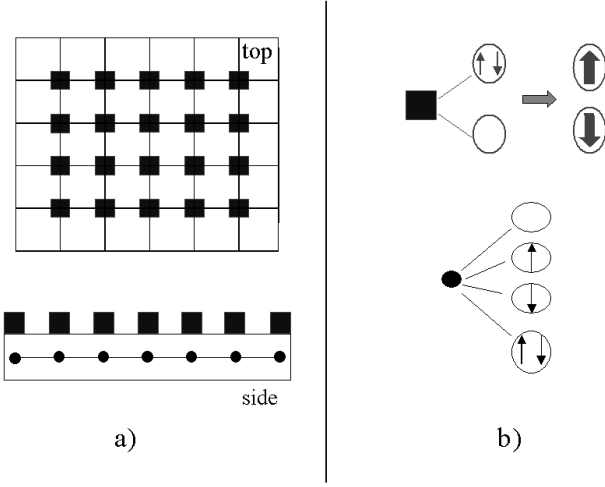


FIG. 1: a) Schematic 1D representation of the BFM on a lattice (top and side view). The bosonic and fermionic particles move on two different arrays having the same periodicity: the fermions are indicated by circles and the bosons by squares on the respective arrays. b) The single site configurations for the pseudospin and fermionic variables.

III. PATH INTEGRAL REPRESENTATION: DERIVATION OF THE EFFECTIVE ACTION

A. Generalities

Let us now construct an effective action which describes the BFM, with the aim to extract the dynamical properties of the low energy degrees of freedom of the phase and amplitude modes for the bosons and fermion-pairs. We start by expressing the partition function in terms of a coherent-state path integral representation⁴⁰, where the fermionic part is formulated by means of the usual Grassmann variables and the bosonic part is described by a pseudospin-coherent state

representation^{40,41}:

$$Z = \prod_i \int D\theta_i D\phi_i D\bar{\Psi}_i D\Psi_i e^{-A[\bar{\Psi}_i, \Psi_i, \theta_i, \phi_i]} \quad (2)$$

where

$$A[\bar{\Psi}_i, \Psi_i, \theta_i, \phi_i] = \int d\tau \sum_i [is(1 - \cos \theta_i) \partial_\tau \phi_i + (\Delta_B - 2\mu) \cos \theta_i] + \sum_{\langle ij \rangle} \bar{\Psi}_j(\tau) G_{ij}^{-1} \Psi_i(\tau). \quad (3)$$

τ denotes the imaginary Matsubara time variable and a Nambu spinor representation for the Grassmann variables, related to that of the original fermionic operators by

$$\bar{\Psi}_i = \begin{pmatrix} c_{i\uparrow} \\ \bar{c}_{i\downarrow} \end{pmatrix} \Psi_i = (\bar{c}_{i\uparrow} c_{i\downarrow}). \quad (4)$$

The pseudospin is described by a bosonic field which in spherical coordinates is given by $\mathbf{s}_i = s(\sin \theta_i \cos \phi_i, \sin \theta_i \sin \phi_i, \cos \theta_i)$ (see Fig. 2). θ_i describes the polar angle of the vector \mathbf{s}_i with respect to the north pole of the z axis, while ϕ_i is the azimuthal angle which defines the angular position of the basal plane projection of this vector. The first term of the action A is the Wess-Zumino term^{41,42}, ensuring the correct quantization of the quantum pseudospin variable. For any path, parameterized by $\phi(\tau)$ and $\theta(\tau)$, the contribution of this term is equal to $i s$ times the surface area of the sphere between this path and the north pole. For closed paths this has exactly the form of the Berry phase⁴³. The second term is linked to the density of the bosons $n_B(\tau)$ through the $\cos \theta(\tau)$ dependence of the pseudospin. Finally, the last contribution of the action contains the coupling between the fermionic and bosonic subsystem through the Green's function G_{ij} , determined by:

$$\begin{pmatrix} K_1 & L \\ L^* & K_2 \end{pmatrix} G_{ij}(\tau - \tau') = \delta(\tau - \tau') \quad (5)$$

where $K_1 = (-\partial_\tau + \mu)\delta_{ij} + t_{ij}$, $K_2 = (-\partial_\tau - \mu)\delta_{ij} - t_{ij}$ and $L = g \sin \theta_i(\tau) e^{i\phi_i(\tau)} \delta_{ij}$ and L^* being the conjugate field of L . Integrating out the fermionic part, one obtains the action in terms of exclusively the bosonic fields:

$$A = -Tr \ln G^{-1} + \int d\tau \sum_i [is(1 - \cos \theta_i(\tau)) \partial_\tau \phi_i(\tau) + (\Delta_B - 2\mu) \cos \theta_i(\tau)] \quad (6)$$

where the trace has to be carried out over all internal as well as space-time indices.

Up to this point no approximation has been made in the derivation of the action which describes the coupling between the bosonic and fermionic degrees of freedom. It is important to notice that the variation of the bosonic variable $\theta(\tau)$ describes both, the density fluctuations of

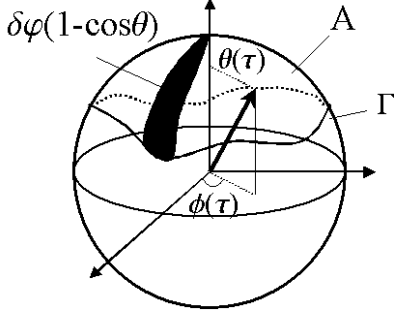


FIG. 2: Spherical representation of the pseudospin \mathbf{s} including the Berry phase factor for one possible trajectory Γ . The Berry phase term is $\exp[i s \oint (1 - \cos[\theta]) \frac{\partial \phi}{\partial \tau}] = e^{i s A}$, where A is the area of the surface enclosed in the trajectory Γ . The black part indicate the differential portion of the surface on the sphere taken respect to the north pole.

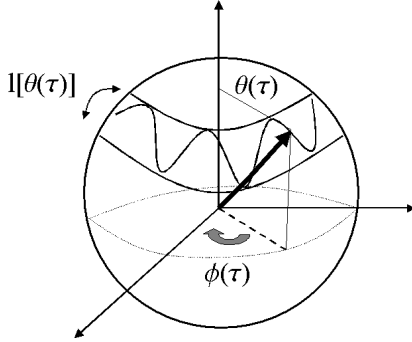


FIG. 3: Representation of a possible path for the local pseudospin motion. The field $l[\theta(\tau)]$ indicates the undulation of the pseudospin vector along the polar direction as it arises from the fluctuations of the average boson density ($\langle \cos \theta \rangle$) and of the pairing amplitude ($\langle \sin \theta \rangle$), while it precesses around the z -axis due to the time evolution of the phase variable $\phi(\tau)$.

the bosonic subsystem (via the projection of the pseudospin vector on the z axis, i.e. $\cos \theta(\tau)$, being the longitudinal component) and the amplitude fluctuations of the fermionic pair field (via the transverse part as projection of the pseudospin vector onto the basal plane, i.e. $\sin \theta(\tau)$) (see Fig. 2)). The variable $\phi(\tau)$ determines the rotational degrees of freedom of the pseudospin vector, expressing its phase dynamics.

For extracting the relevant terms which control the low energy dynamics of the coupled phase and amplitude modes and performing an expansion, which is meaningful in terms of the phase variable, it is judicious to make the following steps: i) gauge away the phase dependence from the term L which permits to separate the trace into a part which does not depend on the phase of the bosonic field and another part that contains only spatial and time variations of $\phi_i(\tau)$ ⁴⁴, ii) rewrite the term L after the gauge transformation, as a sum of two pieces, one not dependent on time (which is linked to the aver-

age density of bosons) and another term containing the fluctuations with respect to its mean value.

The first operation is performed by applying to the operators under the trace the rotation $U_i = e^{i \phi_i(\tau) \sigma_3 / 2}$, where σ_i denote the Pauli matrices. Hence,

$$\text{Tr} \ln G^{-1} = \text{Tr} \ln U \tilde{G}^{-1} U^{-1} = \text{Tr} \ln \tilde{G}^{-1} \quad (7)$$

where

$$\tilde{G}_{ij}^{-1} = \left[-\partial_\tau \sigma_0 + \left[\frac{i}{2} \partial_\tau \phi_i(\tau) - \mu \right] \sigma_3 + g \sin \theta_i(\tau) \sigma_1 \right] \delta_{ij} + t_{ij} e^{i [\phi_i(\tau) - \phi_j(\tau)] \sigma_3 / 2}. \quad (8)$$

Since the bosonic density is fixed in average, we now separate the part which depends on the polar angle in a time independent contribution and its time dependent correction. That is, the term $\sin \theta_i(\tau)$ is decomposed into its average value $\langle \sin \theta_i \rangle$ (which is determined by fixing the density of bosons due to the spherical constraint) plus a time dependent contribution $l[\theta(\tau)]$ which contains the fluctuation around its average value (Fig. 3). Thus this local field, due to the constraint, will describe both: the time dependent variation of the density as well as of the pairing amplitude.

Next, let us write the Green's function in the usual form as

$$\tilde{G}_{ij}^{-1} = G_{0ij}^{-1} + \Sigma_{ij} \quad (9)$$

with $G_{0ij}^{-1} = [-\partial_\tau \sigma_0 - \mu \sigma_3 + \tilde{g} \sigma_1] \delta_{ij}$ and $\Sigma_{ij} = T_{ij} + D_i + K_i$, where

$$D_i = \frac{i}{2} \frac{\partial \phi(\tau)}{\partial \tau} \sigma_3 \quad (10)$$

$$K_i = \tilde{g} l[\theta(\tau)] \sigma_1 \quad (11)$$

$$T_{ij} = t_{ij} e^{-i(\phi_i(\tau) - \phi_j(\tau)) \frac{\sigma_3}{2}} \quad (12)$$

$$\tilde{g} = g \langle \sin \theta \rangle. \quad (13)$$

From this point onward we shall assume the average density of bosons to be homogeneously distributed and thus given by $n_B = \frac{1}{2}(1 + \langle \cos \theta \rangle)$. This implies that the bare coupling g is renormalized to \tilde{g} as a consequence of the spherical constraint of the pseudospin variable. Moreover, since $\langle \sin \theta \rangle = \sqrt{1 - \langle \cos \theta \rangle^2}$ and the density of bosons is fixed in average via a suitable choice of the chemical potential and of the on-site bosonic energy, one can treat the exchange coupling as an external boson-density tunable parameter. The variation in the θ variable is then simply related to the variation of the bosonic density such that if θ varies in the range $[0, \pi]$ then n_B varies in the interval $[0, 1]$.

Before expanding the trace, let us write down explicitly the expression of the zero order Green's function, as it will be frequently used in the following steps:

$$G_{0i}^{\alpha\beta}(\tau) = \frac{1}{\beta} \sum_{\omega_n} G_{0i}^{\alpha\beta}(\omega_n) \exp[-i \omega_n \tau] \quad (14)$$

with

$$G_{0i}^{\alpha\beta}(\omega_n) = \begin{pmatrix} \frac{-i\omega_n + \mu}{\omega_n^2 + \omega_0^2} & \frac{\sqrt{\omega_0^2 - \mu^2}}{\omega_n^2 + \omega_0^2} \\ \frac{\sqrt{\omega_0^2 - \mu^2}}{\omega_n^2 + \omega_0^2} & \frac{-i\omega_n - \mu}{\omega_n^2 + \omega_0^2} \end{pmatrix} \quad (15)$$

and where we introduced $\omega_0 = \tilde{g}$.

B. Second order loop expansion

We now evaluate the contribution of the self energy Σ_{ij} to the effective action. This is done in the usual way by making a loop expansion in the trace⁴⁴. We shall construct that expansion up to second order in the time and space derivatives of the phase variable and in the terms which contain both, the fluctuations of the density and the amplitude. For that purpose we use the standard identity:

$$\begin{aligned} \text{Tr} \ln \tilde{G}^{-1} &= \text{Tr} \ln [G_0^{-1} + \Sigma] \\ &= \text{Tr} \ln G_0^{-1} + \text{Tr} \ln [1 + G_0 \Sigma] \end{aligned} \quad (16)$$

and then expand the second term of this expression up to second order in Σ , such as to keep all the contributions up to quadratic order in the gradient of the phase. This gives

$$\text{Tr} \ln \tilde{G}^{-1} \cong \text{Tr} \ln G_0^{-1} + \text{Tr}[G_0 \Sigma] - \frac{1}{2} \text{Tr}[G_0 \Sigma]^2. \quad (17)$$

The first term of this expression is just a constant and does not contribute to the dynamics. In the second term

$$\text{Tr}[G_0 \Sigma] = \text{Tr}[G_0 T] + \text{Tr}[G_0 D] + \text{Tr}[G_0 K] \quad (18)$$

the only parts different from zero are $\text{Tr}[G_{0i} K_i]$ and $\text{Tr}[G_{0i} D_i]$. $\text{Tr}[G_{0i} T_{ij}]$ gives no contribution once one makes the trace over the site indices. $\text{Tr}[G_{0i} D_i]$ introduces a contribution which is proportional to the chemical potential multiplied by the time derivative of the phase $\sim i\mu \partial_\tau \phi$. We will see below that this contribution describes an effective *off-set charge* (in terms of the terminology of a similar Josephson junction array scenario) and which arises from the total fermionic and bosonic static density distribution via their dependence on the chemical potential. $\text{Tr}[G_{0i} K_i]$ describes the lowest order fluctuations of the bosonic (fermion pair) density due to the presence of the spontaneous pair/hole creation out of the condensate. Its direct evaluation gives a contribution

$$A_1 = \int d\tau_1 d\tau_2 \text{Tr}_{ps} [G_{0i}(\tau_1, \tau_2) K_i(\tau_2, \tau_1)] \quad (19)$$

where Tr_{ps} represents the trace over the internal pseudospin index. With $K_i(\tau_2, \tau_1) = K_i(\tau_1) \delta(\tau_1 - \tau_2)$ and integrating over the Matsubara times gives:

$$\begin{aligned} A_1 &= \int d\tau \text{Tr}_{ps} [G_{0i}(0) K_i(\tau)] \\ &= 2g G_{0i}^{12}(0) \int d\tau l[\theta(\tau)] \\ &= -\frac{g \tanh[\beta\omega_0]}{\omega_0} \int d\tau l[\theta(\tau)], \end{aligned} \quad (20)$$

or in a compact form

$$\begin{aligned} A_1 &= E_1 \int d\tau l[\theta(\tau)] \\ E_1 &= -\frac{g \tanh[\beta\omega_0]}{\omega_0}. \end{aligned} \quad (21)$$

Let us next come to the evaluation of the terms which contribute to the quadratic order in this loop expansion of the trace. The parts which are non zero in these terms are the following:

$$\begin{aligned} A_2 &= \text{Tr} [G_{0i} K_i G_{0i} D_i] \\ A_3 &= \text{Tr} [G_{0i} K_i G_{0i} K_i] \\ A_4 &= \text{Tr} [G_{0i} D_i G_{0i} D_i] \\ A_5 &= \text{Tr} [G_{0i} T_{ij} G_{0j} T_{ji}] \end{aligned} \quad (22)$$

The term A_2 contains processes where the fermionic background locally couples at different times to the fluctuations of the bosonic density and to the phase fluctuations:

$$\begin{aligned} A_2 &= \text{Tr} [G_{0i} K_i G_{0i} D_i] = \\ &\int d\tau_1 d\tau_2 \text{Tr}_{ps} [G_{0i}(\tau_1 - \tau_2) K_i(\tau_2) G_{0i}(\tau_2 - \tau_1) D_i(\tau_1)]. \end{aligned} \quad (23)$$

Since $G_{0i}(\tau_1 - \tau_2)$ depends exclusively on the time differences, we introduce the new variables $\tau = \frac{\tau_1 - \tau_2}{2}$ and $\eta = \frac{\tau_1 + \tau_2}{2}$, after which the integral over the Matsubara time variables becomes:

$$A_2 = \frac{1}{2} \int d\tau d\eta \text{Tr}_{ps} [G_{0i}(\tau) K_i(\eta - \tau) G_{0i}(-\tau) D_i(\tau + \eta)]. \quad (24)$$

As we are interested in the gradient expansion in the bosonic phase and density, related to $\phi(\tau)$ and $\theta(\tau)$, we keep only the lowest order in their time derivatives. This is done by considering the expansion of the product $K_i(\eta - \tau) D_i(\tau + \eta)$ around $\tau = 0$ in order to separate the parts which depend exclusively on the local variable of the Green function and those which are linked to the phase and the density fluctuations. The expansion then reads as follows,

$$\begin{aligned} K_i(\eta - \tau) D_i(\tau + \eta) &\simeq K_i(\eta) D_i(\eta) + \tau \left\{ -\frac{\partial K_i(\eta)}{\partial \eta} D_i(\eta) \right. \\ &\quad \left. + \frac{\partial D_i(\eta)}{\partial \eta} K_i(\eta) \right\} + O(\tau^2) \end{aligned} \quad (25)$$

Due to the symmetry of the Green's functions (even for $\tau \rightarrow -\tau$), the linear contribution of the series expansion of $K_i(\eta - \tau) D_i(\tau + \eta)$ cancels in the effective action after integrating over the time. Hence, to lowest order one obtains,

$$A_2 = \frac{1}{2} \int d\tau d\eta \text{Tr}_{ps} [G_{0i}(\tau) K_i(\eta) G_{0i}(-\tau) D_i(\eta)]$$

$$\begin{aligned}
&= \frac{g}{2} \int d\tau [G_{0i}^{12}(\tau)G_{0i}^{11}(-\tau) + G_{0i}^{11}(\tau)G_{0i}^{12}(-\tau) \\
&- G_{0i}^{22}(\tau)G_{0i}^{12}(-\tau) - G_{0i}^{12}(\tau)G_{0i}^{22}(-\tau)] \times \\
&\int d\eta \frac{i}{2} \frac{\partial \phi(\eta)}{\partial \eta} l[\theta(\eta)]
\end{aligned} \quad (26)$$

or in a compact form,

$$A_2 = E_2 \int d\eta i \frac{\partial \phi(\eta)}{\partial \eta} l[\theta(\eta)] \quad (27)$$

where the coefficient E_2 , after integrating over the Matsubara time, is given by:

$$E_2 = \frac{\tilde{g}^2}{4} \frac{\mu (\beta \omega_0 - \sinh[\beta \omega_0])}{\omega_0^3 (1 + \cosh[\beta \omega_0])}. \quad (28)$$

Let us next consider the term A_3 , which expresses the coupling of the fermionic field to the fluctuations of bosonic density at different times. In order to extract the lowest order gradient contributions, we follow the same procedure as that just used above, giving us:

$$A_3 = \frac{1}{2} \int d\tau d\eta Tr_{ps} [G_{0i}(\tau) K_i(\eta - \tau) G_{0i}(-\tau) K_i(\tau + \eta)] \quad (29)$$

Performing again the expansion in time up to quadratic order in the time derivatives of $K_i(\tau)$, we have:

$$K_i(\eta - \tau) K_i(\eta + \tau) \simeq K_i(\eta)^2 - \tau^2 \frac{\partial K_i(\eta)}{\partial \eta}^2. \quad (30)$$

and hence can rewrite A_3 in the following way:

$$\begin{aligned}
A_3 &= \frac{1}{2} \int d\tau d\eta Tr_{ps} [G_{0i}(\tau) K_i(\eta) G_{0i}(-\tau) K_i(\eta)] + \\
&\frac{1}{2} \int d\tau d\eta Tr_{ps} [G_{0i}(\tau) K_i'(\eta) G_{0i}(-\tau) K_i'(\eta)].
\end{aligned} \quad (31)$$

Carrying out the trace over the internal indices gives:

$$\begin{aligned}
Tr_{ps} [G_{0i}(\tau) K_i(\eta) G_{0i}(-\tau) K_i(\eta)] &= [G_{0i}^{22}(\tau) G_{0i}^{11}(-\tau) \\
+ 2 G_{0i}^{12}(\tau) G_{0i}^{12}(-\tau) + G_{0i}^{11}(\tau) G_{0i}^{22}(-\tau)] &(\tilde{g}^2 l[\theta(\eta)]^2)
\end{aligned} \quad (32)$$

so that

$$A_3 = E_{3a} \int d\eta l[\theta(\eta)]^2 - E_{3b} \int d\eta \frac{\partial l[\theta(\eta)]}{\partial \eta}^2. \quad (33)$$

By evaluating the integrals over the Matsubara times we obtain:

$$\begin{aligned}
E_{3a} &= \frac{\tilde{g}^2}{2} \frac{\beta \omega_0 \tilde{g}^2 + \mu^2 \sinh[\beta \omega_0]}{\omega_0^3 (1 + \cosh[\beta \omega_0])} \\
E_{3b} &= \frac{\tilde{g}^2}{2} \left[\frac{\text{sech}[(\beta \omega_0/2)]^2 (-6\beta \mu^2 \omega_0 + \beta^3 (-\mu^2 \omega_0^3 + \omega_0^5))}{24\omega_0^5} \right. \\
&\quad \left. + \frac{6\mu^2 \sinh[\beta \omega_0]}{24\omega_0^5} \right].
\end{aligned} \quad (34)$$

Applying the same procedure for the evaluation of A_4 (which has the same functional form as A_3 but involving the coupling between the fermionic degrees of freedom and the phase velocity at different time) one gets an analogous expression, providing we discard terms of higher order in the derivatives than $[\frac{\partial \phi(\tau)}{\partial \tau}]^2$:

$$\begin{aligned}
A_4 &\simeq \frac{1}{2} \int d\tau d\eta Tr_{ps} [G_{0i}(\tau) D_i(\eta) G_{0i}(-\tau) D_i(\eta)] \\
&= \frac{1}{8} C_0 \int d\eta \left(\frac{\partial \phi(\eta)}{\partial \eta} \right)^2.
\end{aligned} \quad (35)$$

Here

$$C_0 = \int d\tau \{ -[G_{0i}^{11}(\tau) G_{0i}^{11}(-\tau) - 2 G_{0i}^{12}(\tau) G_{0i}^{12}(-\tau) + G_{0i}^{22}(\tau) G_{0i}^{22}(-\tau)] \}, \quad (36)$$

which after evaluating the integration over the Matsubara times gives,

$$C_0 = \frac{\text{sech}[(\beta \omega_0/2)]^2 (\beta \mu^2 \omega_0 + \tilde{g}^2 \sinh[\beta \omega_0])}{2\omega_0^3}. \quad (37)$$

Finally, we come to the evaluation of the last term A_5 which involves the coupling between phase fluctuations on different sites and the process of single particle hopping. This contribution will yield terms which are quadratic in the time derivative of the phase (*charging like*, in the terminology of a similar Josepshon junction array scenario) and moreover will generate an effective hopping induced inter-site phase coupling which turns out to be similar to the Josepshon coupling in arrays of superconducting grains.

Considering again the lowest order gradient expansion contributions we have:

$$\begin{aligned}
A_5 &= Tr [G_{0i} T_{ij} G_{0j} T_{ji}] \\
&\simeq \frac{1}{2} \int d\tau d\eta Tr_{ps} \{ G_{0i}(\tau) T_{ij}(\eta) G_{0j}(-\tau) T_{ji}(\eta) - \\
&\tau^2 \left[G_{0i}(\tau) \frac{\partial T_{ij}(\eta)}{\partial \eta} G_{0j}(-\tau) \frac{\partial T_{ji}(\eta)}{\partial \eta} \right] \} \\
&= P_1 + P_2
\end{aligned} \quad (38)$$

The terms $P_{1,2}$ are given by:

$$\begin{aligned}
P_1 &= -t_{ij}^2 \int d\tau [G_{0i}^{12}(\tau) G_{0i}^{12}(-\tau)] \int d\eta \cos[\phi_i(\eta) - \phi_j(\eta)] \\
P_2 &= -\frac{t_{ij}^2}{8} \int d\tau \tau^2 [G_{0i}^{11}(\tau) G_{0j}^{11}(-\tau) + G_{0i}^{22}(\tau) G_{0j}^{22}(-\tau)] \\
&\int d\eta \left[\frac{\partial \phi_i(\eta)}{\partial \eta} - \frac{\partial \phi_j(\eta)}{\partial \eta} \right]^2.
\end{aligned} \quad (39)$$

By carrying out the integration over the Matsubara times one obtains for the hopping induced inter-site amplitude

and phase coupling:

$$\begin{aligned} P_1 &= E_J \int d\eta \cos[\phi_i(\eta) - \phi_j(\eta)] \\ E_J &= \frac{t_{ij}^2 \tilde{g}^2 (-\beta\omega_0 \text{sech}[\beta\omega_0/2]^2 + 2 \tanh[\beta\omega_0/2])}{8\omega_0^3}. \end{aligned} \quad (40)$$

Similarly, by calculating the coefficient of the term P_2 , one obtains the strength of a *mutual capacitance* term (in the terminology of a similar Josephson junction array scenario) for neighboring effective sites:

$$\begin{aligned} P_2 &= \frac{1}{8} C_1 \int d\eta \left[\frac{\partial \phi_i(\eta)}{\partial \eta} - \frac{\partial \phi_j(\eta)}{\partial \eta} \right]^2 \\ C_1 &= -t_{ij}^2 \int_{-\beta/2}^{\beta/2} d\tau \tau^2 [G_i^{11}(\tau) G_i^{11}(-\tau) + G_i^{22}(\tau) G_i^{22}(-\tau)] \\ &= -t_{ij}^2 \left[\frac{\exp[\beta\omega_0] (-6\beta\tilde{g}^2\omega_0 + \beta^3\omega_0^3(\mu^2 + \omega_0^2))}{12(1 + \exp[\beta\omega_0])^2\omega_0^5} \right. \\ &\quad \left. + \frac{6\tilde{g}^2 \sinh[\beta\omega_0]}{12(1 + \exp[\beta\omega_0])^2\omega_0^5} \right]. \end{aligned} \quad (41)$$

The final effective action is then given by the sum of all the terms evaluated above which are grouped together in form of three different contributions:

$$S = \int d\tau [S_\phi + S_{\phi-\theta} + S_\theta]. \quad (42)$$

S_ϕ , S_θ , $S_{\phi-\theta}$ are the contributions arising from exclusively (i) the phase dynamics, (ii) the fluctuations of the bosonic density and of the amplitude $l[\theta]$ and (iii) the coupling between them. They are given by:

$$\begin{aligned} S_\phi &= \sum_i \frac{1}{8} (C_0 + zC_1) \left(\frac{\partial \phi_i(\tau)}{\partial \tau} \right)^2 \\ &\quad - \sum_{(i,j)} \left[\frac{1}{8} C_1 \frac{\partial \phi_i(\tau)}{\partial \tau} \frac{\partial \phi_j(\tau)}{\partial \tau} + E_J \cos[\phi_i(\tau) - \phi_j(\tau)] \right] \\ &\quad + i \sum_i \mu \frac{\partial \phi_i(\tau)}{\partial \tau} \\ S_\theta &= \sum_i \left[E_1 l[\theta_i(\tau)] + E_{3a} l[\theta_i(\tau)]^2 - E_{3b} \left[\frac{\partial l[\theta_i(\tau)]}{\partial \tau} \right]^2 \right] \\ S_{\phi-\theta} &= \sum_i \left[i s \frac{\partial \phi_i(\tau)}{\partial \tau} (1 - \cos[\theta_i(\tau)]) \right. \\ &\quad \left. + i E_2 \frac{\partial \phi_i(\tau)}{\partial \tau} l[\theta_i(\tau)] \right]. \end{aligned} \quad (43)$$

We expand the Berry phase contribution (the first term in $S_{\phi-\theta}$) up to second order in $l[\theta(\tau)]$ and subsequently redefine this field as $l[\theta(\tau)] = a + b\tilde{l}[\theta(\tau)]$, with the time independent constants a, b chosen in such a way as to eliminate any terms linear in \tilde{l} in the the action S . We then find two contributions to the action which are linear in $\partial \phi(\tau)/\partial \tau$: One which is time independent and which

can be absorbed into the chemical potential and another one which is quadratic in $\tilde{l}[\theta(\tau)]$. With this, $S_{\phi-\theta}$ can be rewritten as:

$$S_{\phi-\theta} = \sum_i i \frac{\partial \phi_i(\tau)}{\partial \tau} \bar{q}_i(\tau), \quad (44)$$

where $\bar{q}_i(\tau)$ is quadratic in $\tilde{l}[\theta(\tau)]$.

The effective action thus constructed for the BFM (Eqs. (42,43)) is similar to that of Josephson junction arrays with nearest neighbor, as well as, self *capacitance*, *Josephson* coupling and *off-set charge* terms. The action for the BFM goes however beyond that for such Josephson junction arrays in the following respects. We have extra terms which control the dynamics of the amplitude modes, given by S_θ and an intrinsic Berry phase term which gives rise to a direct phase - amplitude coupling (Eq. 44), where the dynamical amplitude fluctuations would correspond to a time dependent *offset charge* term in an analogous Josephson junction array picture. Finally, the Berry phase term, being an intrinsic topological term, may give rise to a Magnus force on a vortex, as it will be discussed in section V.

IV. THE SUPERCONDUCTOR - INSULATOR PHASE BOUNDARY

In a first approach to analyze the stability region of long-range phase superconducting coherence, we examine the effective action (eqs. 42,43), restricting ourselves to the phase only part of it. This means a study of the SIT driven by a competition between the phase coherence induced by pair hopping and the disrupting effect of local boson density (or equivalently pair field amplitude) fluctuations in the presence of a source term for the bosons which controls the global boson density via an effective chemical potential. Within such an approximation our study is equivalent to that of Josephson junction arrays, except that the effective coupling constants entering in such an action depend in a highly non trivial way on the parameters which characterize the original BFM Hamiltonian. This, as we shall see, will lead to novel features concerning the phase diagram with a SIT for the BFM when we examine it in terms of the boson-fermion exchange coupling g/t and the boson concentration n_B .

We shall determine the phase diagram by means of the so-called coarse graining approximation which has been successfully applied for this kind of problem and which permits to capture the relevant qualitative and quantitative features of such a Josephson junction array like action^{45,46}. It is known that, as a consequence of the uncertainty relation between the phase ϕ_i and the pair number operator $Q_i = i \frac{\partial}{\partial \phi_i}$, the system can switch from a phase ordered to a disordered state. An essential part of this study will concern how the relevant parameters of the BFM Hamiltonian influence the equivalent amplitudes of the Josephson coupling, the *capacitance* and

the *off-set charge* terms. Thus, at zero temperature, by fixing the bosonic density distribution, we find that the variation of the ratio between the Josephson coupling energy E_J and the *charging* energy $E_C = C_0^{-1}/2$, which controls the phase-density interplay, increases from zero, goes through a maximum and then decreases to zero with increasing g/t . Or else, if one fixes the coupling g/t , then by varying the density of the bosonic distribution one is able to control the effective coupling which appears in E_J and C_{ij} by varying the average bosonic density $\tilde{g} = 2g\sqrt{n_B(n_B - 1)}$. It is thus immediately evident that in the BFM scenario there is a non-trivial interplay between the renormalization of the Josephson coupling and the charging effect. If one goes to the limit of empty ($n_B = 0$) and full bosonic occupation ($n_B = 1$), the effective coupling $\tilde{g} \rightarrow 0$, and the critical temperature consequently reduces to zero.

The general form of the action we then have to examine is given by:

$$S_{\text{phase}} = \int_0^\beta \sum_i \frac{i}{2} \frac{\partial \phi_i(\tau)}{\partial \tau} q_i - E_J \sum_{ij} \alpha_{ij} \cos[\phi_i(\tau) - \phi_j(\tau)] + \sum_{i,j} \frac{1}{8} \frac{\partial \phi_i(\tau)}{\partial \tau} C_{ij} \frac{\partial \phi_j(\tau)}{\partial \tau}. \quad (45)$$

The first term describes the effect of a static *off-set charge* (q_i). The second term contains the physics of the pair hopping processes, an analog to the Josephson tunnelling processes, and has a coupling strength E_J with $\alpha_{ij} = 1$ if (i, j) are nearest neighbors and zero otherwise. Finally, the third term describes the *charging* term arising from the local exchange of boson and fermion pairs and from quasi-particle hopping between nearest neighbor sites. The strength of this *charging* type interaction is given by $C_{ij} = (C_0 + zC_1)\delta_{i,j} - C_1 \sum_p \delta_{i,i+p}$ which represents an effective general *capacitance* matrix, with the vector p running over the nearest neighbors. C_0 denotes the *self-capacitance* and C_1 the mutual one (z being the coordination number).

As mentioned before, to extract the phase diagram we make use of the coarse graining approximation⁴⁵. The main idea of this approach is to introduce a Hubbard-Stratonovich auxiliary field which is conjugate to the average of $\langle e^{i\phi_i} \rangle$ and which plays the role of an order parameter for the transition from a superconducting to an insulating state. Since the phase transition has a continuous character, one can expand the action in powers of the auxiliary field and determine the occurrence of phase coherence by looking at the coefficients of the quadratic term in the limit of long-wavelengths and zero frequency.

We briefly sketch the main steps of such an approximation and adapt it to the present scenario of the BFM. The partition function for S_{phase} is given by the sum of all the possible paths of the phase variables in the imag-

inary time and in the real space:

$$Z = \int \prod_i D\phi_i \exp \left\{ \int_0^\beta \sum_i -\frac{i}{2} \frac{\partial \phi_i(\tau)}{\partial \tau} q_i(\tau) + E_J \sum_{ij} \alpha_{ij} \cos[\phi_i(\tau) - \phi_j(\tau)] - \frac{1}{8} \sum_{ij} \frac{\partial \phi_i(\tau)}{\partial \tau} C_{ij} \frac{\partial \phi_j(\tau)}{\partial \tau} \right\}$$

To perform the Hubbard-Stratonovich transformation, one rewrites the *Josephson coupling* term as $\frac{E_J}{2} \sum_{ij} \exp[i\phi_i] \alpha_{ij} \exp[i\phi_j]$ and then, by using the usual Gaussian identity, introduces an auxiliary field $\psi_i(\tau)$. The partition function Z then becomes:

$$Z = \int \prod_i D\psi_i^* D\psi_i D\phi_i \exp \left\{ \int_0^\beta d\tau \left[\left(\frac{-2}{E_J} \right) \sum_{ij} \psi_i^* \alpha_{ij}^{-1} \psi_j \right] \exp \left\{ \int_0^\beta d\tau \left[\sum_i \frac{i}{2} \frac{\partial \phi_i(\tau)}{\partial \tau} q_i(\tau) - \sum_i (\psi_i e^{i\phi_i} - \psi_i^* e^{-i\phi_i}) - \sum_{i,j} \frac{1}{8} \frac{\partial \phi_i(\tau)}{\partial \tau} C_{ij} \frac{\partial \phi_j(\tau)}{\partial \tau} \right] \right\} \right\}. \quad (46)$$

Hence, starting from the effective action for the auxiliary field, one can perform an expansion up to second order in ψ and thus derive the usual Ginzburg Landau type free energy functional which permits to determine the boundary line between the superconducting and the insulating state. The corresponding effective action for these auxiliary fields ψ, ψ^*

$$S_\psi = \ln \left[\int \prod_i D\phi_i \exp \left(\int_0^\beta d\tau \left\{ \sum_i \frac{i}{2} \frac{\partial \phi_i(\tau)}{\partial \tau} q_i(\tau) - \sum_i (\psi_i(\tau) e^{i\phi_i(\tau)} - \psi_i^*(\tau) e^{-i\phi_i(\tau)}) - \sum_{i,j} \frac{1}{8} \frac{\partial \phi_i(\tau)}{\partial \tau} C_{ij} \frac{\partial \phi_j(\tau)}{\partial \tau} \right\} \right) \right] \quad (47)$$

is then expanded to second order, giving:

$$S_\psi = \int_0^\beta d\tau d\tau' \chi_{ij}(\tau, \tau') \psi_i^*(\tau) \psi_j(\tau) + O(\psi^4), \quad (48)$$

where

$$\chi_{ij}(\tau, \tau') = \langle e^{i[\phi_i(\tau) - \phi_j(\tau')]} \rangle_0 \quad (49)$$

denotes the two-time phase correlator, which is equal to the second derivative with respect to the auxiliary field ψ and its conjugate at different time and space positions, evaluated around their zero values. By performing the functional derivation, one gets the following expression for it:

$$\chi_{ij}(\tau, \tau') = \frac{\int \prod_i D\phi_i e^{i[\phi_i(\tau) - \phi_j(\tau')]} \exp[-S_0]}{\int \prod_i D\phi_i \exp[-S_0]} \quad (50)$$

with S_0 being the part of the action which contains exclusively the *charging* contributions, i.e. :

$$S_0 = \exp \left[\int_0^\beta d\tau \sum_i \frac{i}{2} \frac{\partial \phi_i(\tau)}{\partial \tau} q_i(\tau) - \sum_{i,j} \frac{1}{8} \frac{\partial \phi_i(\tau)}{\partial \tau} C_{ij} \frac{\partial \phi_j(\tau)}{\partial \tau} \right]. \quad (51)$$

According to the scheme outlined above, we now develop the partition function Z up to second order in the auxiliary fields, thus putting it into a familiar form:

$$Z = \int \prod_i D\psi_i^* D\psi_i e^{-F_\psi} \quad (52)$$

with

$$F_\psi = \int_0^\beta d\tau d\tau' \sum_{ij} \psi_i^*(\tau) \left[\alpha_{ij} \delta(\tau - \tau') - \chi_{ij}(\tau, \tau') \right] \psi_j(\tau'). \quad (53)$$

The determination of the conditions for the boundary line between the superconducting and the insulating phase then reduces to the explicit evaluation of $\chi_{ij}(\tau, \tau')$. The result for that has been first obtained in Ref. [48], and we sketch below the main steps of this derivation.

In determining $\chi_{ij}(\tau, \tau')$ it has turned out to be essential to treat the phase variable ϕ_i in a compact form. In order to separate the imaginary time evolution of the phase into a periodic part $\bar{\phi}_i(\tau)$ and a non-periodic part, one introduces the following parameterization:

$$\phi_i(\tau) = \bar{\phi}_i(\tau) + \frac{2\pi n_i \tau}{\beta} \quad (54)$$

with $\bar{\phi}_i(0) = \bar{\phi}_i(\beta)$ and n_i being an integer which counts how many times the phase winds over an angle which is a multiple of 2π .

The use of such a relation allows to express the sum over all ϕ_i as an integration over $\bar{\phi}_i$ plus a sum over all the possible integer values of the winding number n_i . After performing a number of suitable algebraic operations⁴⁶, one ends up with the following expression for the two-time phase correlator:

$$\chi_{ij}(\tau) = \delta_{ij} \frac{\exp[-2C_{ii}^{-1}|\tau|]}{\sum_{\{n_i\}} \exp \left[- \sum_{ij} 2\beta C_{ij}^{-1} N_i N_j \right]} \times \sum_{\{n_i\}} \exp \left[- \sum_{ij} 2\beta C_{ij}^{-1} N_i N_j - \sum_k 4C_{ik}^{-1} N_k \tau \right], \quad (55)$$

where $N_i = q_i/2 + n_i$. This two-time phase correlator is local in space and its time dependence follows an exponential behavior, if one assumes that the static *offset charge* q_i has a distribution which is homogeneous

in space. In the general case, of a dynamic *offset charge* $q_i(\tau)$, this time dependence will be modified and can lead to qualitative changes in the nature of the SIT.

Having obtained the expression for the local two-time phase correlator in the time and space representation, we can now express the effective Ginzburg-Landau free energy functional in the Fourier space in the following way:

$$F_\psi = \frac{1}{\beta L} \sum_{n,k} \psi_k^*(\omega_n) [\alpha_k^{-1} - \chi_k(\omega_n)] \psi_k(\omega_n). \quad (56)$$

Expanding the inverse matrix α_{ij}^{-1} in the form $\alpha_k^{-1} = (1/z) + k^2(a^2/z^2) + \dots$, this free-energy functional finally is written as:

$$F_\psi = \frac{1}{\beta L} \sum_{k,n} \left[\frac{2}{zE_J} - \chi_0 + ak^2 + b\omega_n^2 + \dots \right] |\psi_k(\omega_n)|^2. \quad (57)$$

The transition line is then given by the condition that the coefficient of the quadratic term vanishes in the limit of vanishing k and ω , that is:

$$1 - \frac{zE_J}{2} \chi_0(0) = 0. \quad (58)$$

For a quantitative analysis, one has to determine the explicit zero frequency limit of the two-time phase correlator. As mentioned above, the inclusion of the time dependent *offset-charge* coming from the fluctuations of the bosonic density can modify the low frequency behavior of the local phase correlations.

In order to get a first insight into the underlying physics at play here we evaluate the two time phase correlator in the limit of a purely local *capacitance* (the so called *self-charging* limit), by keeping only the on-site part in the original structure of the *capacitance* matrix. Under those conditions the evaluation of the two time phase correlator at finite frequency gives:

$$\chi_{ii}(\omega_n) = \frac{1}{Z_0} \sum_{\{n_i\}} F[n_i] \left(\frac{1}{C_{ii}^{-1}} - [i\omega_n - 4C_{ii}^{-1}n_i] \right) \quad (59)$$

with $F[n_i] = \exp[-\sum_i 2\beta C_{ii}^{-1}n_i^2]$ and $Z_0 = \sum_{\{n_i\}} F[n_i]$.

With these expressions we finally can cast Eq. 58 in the form

$$1 - \frac{zE_J}{4E_C} \frac{\sum_n \frac{\exp[-4\beta E_C n^2]}{1-4n^2}}{\sum_m \exp[-4E_C \beta m^2]} = 0, \quad (60)$$

expressed in terms of the *Josephson coupling* E_J and the *charging energy* $E_C = C_{ii}^{-1}/2$, as given by the Eqs.(37,40,41).

Given this defining equation for the boundary between the superconducting and the insulating phase, we want

to see now how the intrinsic dependence of the Josephson and charging energy on the exchange coupling and the density of bosons, manifests itself in the competition between phase and boson density degrees of freedom. In the loop expansion given in the previous section for the BFM, we have obtained the amplitude of the intersite phase coupling and the *charging* effect as a function of temperature, the ratio g/t and the bosonic density. It is the non trivial dependence of those coupling constants on the parameters of the original microscopic Hamiltonian of the BFM which leads to the interesting features in the competition between the phase and boson density fluctuations. Thus, upon increasing the pair exchange coupling it is expected that the *Josephson* coupling is reduced while the *charging* energy is increased due to the local transfer between fermions pairs and bosons as well as due to the single particle hopping of fermions between nearest neighbor sites.

We determine the critical line for two different cases in order to highlight the role played by, on the one hand, the coupling and, on the other hand, by the effective boson density. In Fig. 4, we illustrate the transition line T_ϕ , separating a phase coherent state from a phase disordered one, as a function of the coupling strength and the effective boson density n_B . The evolution of the transition line is non monotonic as a function of g/t and goes through a maximum at $g_{max} \sim t$. The quantum critical point, where the SIT occurs, is given by $g_{crit} \sim 2t$. The critical behavior close to the transition is that of an XY model in $d+1$ dimensions.

More interesting still is the behavior of T_ϕ as a function of the n_B . The variation of T_ϕ is qualitatively different for the different parameter regimes: a) the weak coupling case for $g < g_{max}$, b) the intermediate one with $g_{max} < g < g_{crit}$ and c) the strong coupling limit for $g > g_{crit}$ (see Fig. 5). Going from the limit *a* to *c* we find that

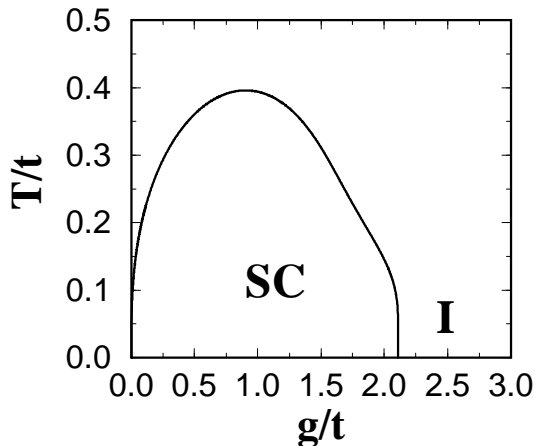


FIG. 4: Phase diagram as a function of g/t describing the boundary between an insulating(I) and superconducting(SC) phase with long range phase coherence for the case of fully symmetric limit ($\Delta_B = 0, n_{F\uparrow} = n_{F\downarrow} = 2n_B = 1$).

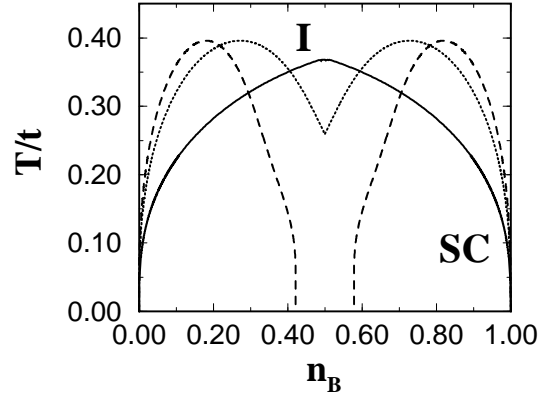


FIG. 5: Phase boundary lines between the superconducting(SC) and insulating(I) state at different values of the coupling and upon varying the average density of bosons. Solid, dotted and dashed line stand for values of the coupling constants for which $g < g_{max}$, $g_{crit} > g > g_{max}$ and $g > g_{crit}$, respectively.

with increasing n_B the critical line decreases to zero, goes through a maximum starting at a finite $T^\phi(n_B = 1)$ and finally shows a SIT at a critical density.

V. ROLE OF BERRY PHASE TERM: AN INTRINSIC MAGNUS FORCE ON VORTICES

In the preceding section we have analyzed the basic physics resulting from the effective action by neglecting, (i) the influence of any feedback between the density and amplitude fluctuations (included in the field $l[\theta_i]$) on the phase dynamics, and (ii) the contribution from the Berry phase term, responsible for the correct quantization of the pseudospin variable, which is given by the integral over all the possible paths of $is(1 - \cos[\theta_i])\partial_\tau\phi_i$. In this section, we discuss the consequences of the presence of such a Berry phase term in the case where the phase action has a vortex solution. The existence of a vortex solution is assured for 2D systems where the phase correlations are described by a XY type dynamics. We shall show here that the Berry phase term will produce an intrinsic Magnus Force on the vortex which is analogous to the Lorentz force for a charged particle, whose effective magnetic field depends on the spatial distribution of the θ field. Using the correspondence with the bosonic pseudospin variable, this implies a relation with the spatial density distribution of the bosons. The Magnus force has been widely discussed in the context of normal BCS type superconductors where it has been shown that it arises from the Berry phase caused by the adiabatic motion of a vortex along a closed loop coming back to its starting position⁴⁸. The adiabatic vortex motion on a loop in the superconducting state turns out to be affected by an effective magnetic field generated by the supercurrent arising from the gradient of the phase which encircles such a

vortex. In the BFM scenario discussed here, we find that on top of the usual contribution due to the superfluid electrons, there is an intrinsic Berry phase term which will generate such a Magnus force and which has an intensity proportional to $2n_B - 1$, where n_B is the bosonic density. This Magnus force, arising solely from the quantum nature of the pseudospin variable, is clearly independent on any superconducting state, and does not require a coherent superfluid current induced by the presence of the vortex itself. Moreover, since n_B can be controlled externally, one has the possibility to tune the strength of the effective magnetic field acting on the vortex, and hence to alter its dynamical properties. This will result in a possible measurable effect on transport coefficients, such as the Hall coefficient, resistance, etc.

To be more explicit and following the procedures used in different approaches treating with the Magnus force problem, let us assume that one has a vortex centered at the position R . Let us furthermore consider that we are in the continuum and at zero temperature so that the Berry phase term in the Lagrangian now reads $L_B = \int d^2r \cos[\theta(r, t)] \partial_t \phi(r, t)$. The contributions linear in $\partial_t \phi(r, t)$ and having time independent coefficients do not contribute to the dynamics. Moreover, since we consider that the phase ϕ and the variable θ are linked to a vortex solution centering at R , it is judicious to introduce a new relative variable $r - R$, which defines the intrinsic position dependence of the variables with respect to the vortex center.

Starting from the above given expression for the Lagrangian L_B and making a simple change in the time derivative we can rewrite it in the following form:

$$\begin{aligned} L_B &= \int d^2r \cos[\theta(r - R, t)] \partial_t \phi(r - R, t) \\ &= \frac{d\mathbf{R}}{dt} \int d^2r \cos[\theta(r - R, t)] \nabla_r \phi(r - R, t). \end{aligned} \quad (61) \quad (62)$$

At this point, one recognizes that the term $\int d^2r \cos[\theta(r - R, t)] \nabla_r \phi(r - R, t)$ plays the role of an effective vector potential \mathbf{A}_B and $\frac{d\mathbf{R}}{dt}$ represents the velocity of the vortex. This Lagrangian is thus equivalent to that of a charged particle in presence of an effective magnetic field in the z-direction given by $\mathbf{B} = \nabla \times \mathbf{A}$. \mathbf{B} hence creates a Magnus force $\mathbf{F}_M = -(\frac{d\mathbf{R}}{dt} \times \hat{\mathbf{z}}) 2\pi \cos[\theta_0]$ which is transverse and proportional to the vortex velocity and whose strength is linked to the magnitude of this effective magnetic field $B_z = -2\pi \cos[\theta_0]$. Its magnitude is given by the asymptotic value of the background bosonic density at large distance $\cos[\theta_0]$ for $r \rightarrow \infty$.

In case of a 2D superconductor described by the BFM scenario, the usual effective magnetic field $-2\pi\rho_s$ arising from the supercurrent of the electrons circling the vortex has thus to be supplemented by the above mentioned contribution B (intrinsic to such a scenario) and permits in principle to modulate the total Magnus force upon changing the density of the bosons. Clearly, a full description of such an eventuality, which is beyond the scope of the present analysis, has to be studied in more

detail and has to include the derivation of the effective action in presence of magnetic field such as to determine the superfluid density of the electrons ρ_s .

At this stage we simply want to point out that the sign of the magnetic field arising from the Berry phase term in the action of the BFM changes if one tunes the bosonic density between zero and unity. In terms of the variable $\cos[\theta]$ this implies a variation in the range $[-1, 1]$. In other words, the sign of the intrinsic Berry phase induced Magnus force will change at $n_B = 1/2$ when going from the limit of small density of bosons to high density. This extra contribution to the Magnus force has to be added of course to the conventional one known for standard superconductors.

VI. COMPARISON WITH NEGATIVE-U AND BOSE-HUBBARD MODELS

A frequently asked question concerns the qualitative differences which exist between the BFM and similar scenarios such as the negative-U Hubbard model and pure bosonic systems such as the Bose-Hubbard model and Josephson junction arrays.

Let us start with a comparison of the BFM and the negative-U Hubbard model which has been studied in a great variety of different approaches and discussed especially in connection with the BCS-BEC crossover. The main issue of such a comparison is the occurrence or not of a quantum critical point in the negative-U Hubbard scenario at a finite value of the coupling U which introduces the pairing among the electrons. As we have shown in this paper, in the BFM there exists a critical value for the exchange coupling g/t responsible of the pairing in the fermionic subsystem above which the system undergoes a SIT. We have seen that within a path integral formalism, one can extract an effective action describing such a transition as being primarily due to the interplay between the local quantum pair exchange (between a boson and a fermion pair) and the non-local inter-site pair hopping leading to a phase dynamics originating from the itinerancy of the fermions. Due to the non-trivial dependence of the strength of these two competing mechanisms on the microscopic parameters of the BFM, there occurs a SIT for a finite density as well as finite exchange coupling g/t . For the negative-U Hubbard scenario this is not the case and the SIT does not occur at any finite value of the ratio U/t ⁴⁹. This model merely describes a continuous cross-over between a BCS superconductor and a BEC of tightly bound electron pairs as U is increased from 0 to ∞ .

In order to better understand this difference between the BFM and the negative-U Hubbard model, let us consider two scenarios in equivalent situations, i.e., the half filled band case for the negative-U Hubbard model and the fully symmetric case for the BFM (with the bosonic level lying in the middle of the fermionic band such that both, the fermionic band as well as the bosonic level are

half occupied). We then address the question how, in the strong coupling limit, the pair hopping is generated out of the basic configurations of states and processes which contribute in this regime.

For the case of the negative- U Hubbard model with fermions interacting via a local attractive potential, the ground state wave function in the limit of $U \rightarrow \infty$ and at half-filling is highly degenerate and is composed of all the possible configurations comprising equal distribution of zero and doubly occupied site. In this limit, one can perform a mapping of this model on the hard-core boson model described by the Hamiltonian:

$$H_U = - \sum_{\langle ij \rangle} \frac{2t_{ij}^2}{U} b_i^+ b_j + \sum_{\langle ij \rangle} \frac{2t_{ij}^2}{U} n_i n_j - \tilde{\mu} \sum_i n_i. \quad (63)$$

b_i^+ (b_i) and $n_i = b_i^+ b_i$ stand for the creation (destruction) operators of hard-core bosonic particles (tightly bound electron pairs) and for their density operator, respectively. Due to the presence of a coupling, which is isotropic both in the boson hopping and in the charge interacting channel, one has a superconducting state for any finite value of the ratio U/t , and possibly a super-solid phase due to the symmetry of the charging interactions characterized by a coexistence of diagonal and off-diagonal long range order. This implies that the quantum critical point is strictly pushed to $U \rightarrow \infty$.

In the BFM the phase space in the large g/t limit is completely different. In the fully symmetric case of this model (corresponding to a total density equal twice the number of sites ($n_{tot} = 2$)), the ground state configuration is given by a wave function where all the bosons are strongly coupled with pairs of fermions, thus resulting in a product state of local bonding states:

$$|\psi_0\rangle = \prod_i \frac{1}{\sqrt{2}} (\rho_i^+ + \tau_i^+) |0\rangle. \quad (64)$$

This wave function can be viewed as a ferromagnetic Ising-type state in the sense that it is made up of a ferro-type bonding order and is not degenerate with respect to all the other possible local configurations. Moreover, by construction it contains bonding-bonding correlations on a long range (bonding solid), but has zero phase correlation length. The latter can be shown by evaluating the static correlation function for the bosons or, equivalently, the fermion pairs which gives $\langle \psi_0 | \rho_i^+ \rho_j | \psi_0 \rangle = 0$ and $\langle \psi_0 | \tau_i^+ \tau_j | \psi_0 \rangle = 0$ for any distance $|i - j|$.

Next, let us consider the low energy configurations which are mixed into such a $g/t = \infty$ ground state when the kinetic energy operator $H_t = \sum_{i \neq j, \sigma} t_{ij} (c_{i\sigma}^+ c_{j\sigma} + H.c.)$ is switched on. Applying H_t to $|\psi_0\rangle$, there will occur states which are separated by an energy gap with respect to $|\psi_0\rangle$. The low energy states to be considered as the relevant quantum mixed configurations are of non-bonding nature, such as $C_{i\sigma}^+ |0\rangle = c_{i\sigma}^+ |0\rangle$ and $S_{i\sigma}^+ |0\rangle = \rho_i^+ c_{i\sigma}^+ |0\rangle$. In order to construct those states, let us to begin with consider a local excitation on two adjacent sites i, j given by $S_{i\sigma}^+ C_{j-\sigma}^+ |0\rangle$ in the background

of bonding states. This configuration can be seen as a ferro-type order interrupted by two domain walls of non-bonding/bonding nature. Now let us consider the dynamics of such objects and under which conditions they can be rendered itinerant in a way which leaves the number of bonding configurations unchanged. The problem is thus analogous to that of domain wall dynamics in an Ising type systems. It so happens (a detailed discussion of this is beyond the scope of the present study and will be given at a later stage elsewhere), that the local degeneracy of the non-bonding states will be removed by the action of the kinetic term. In this way it induces a global lowering of their excitation energy due to the dispersion of each domain wall which is of the form

$$E_{DW}(k) = \frac{1}{2} [\epsilon_k + 6g \pm \sqrt{(\epsilon_k - 2g)^2 + 4t^2}] \quad (65)$$

and where $\epsilon(k) = -2zt \cos k$. Only after the condensation of the domain wall like excitations in the presence of the bonding state background one can meet the conditions for setting up long range boson and fermion-pair phase correlations.

This simple sketch of the nature of the excitations in the BFM implies that the onset of phase coherence cannot be activated for an arbitrarily small hopping amplitude since one has to overcome the energy gap between the ground state and the manifold of non-bonding states, which is of the order of $\sim 2g$. As we can see from the expression for the dispersion of the domain walls, this is achieved when $E_{DW}(k)$ becomes zero for $k = 0$, i.e., when $g/t = \frac{1}{4}(2 + \sqrt{2 + 4z^2})$. For $z=2$ this gives $g/t \sim 2.0$, which determines when those defects-like excitations become gapless and thus induce a proliferation within the background of the bonding states. It is worth pointing out that this value for the critical exchange coupling matches rather well the value obtained from our functional integral approach discussed in this paper (see Fig 3).

The picture which hence emerges for the BFM is very similar to that of a XY model in a transverse field. There, the hopping term is responsible of the XY dynamics, while the local boson-fermion pair exchange provides the role of the effective transverse field. This itself is already a strong indication that this model is of different nature to that of the negative- U Hubbard scenario, for which we know that it is, on the contrary, akin to an isotropic Heisenberg model.

A further essential difference between the BFM and the negative- U Hubbard model can be seen from their respective path-integral formulations. When the path integral representation is considered for this model⁵⁰, the first step is to make use of a Hubbard-Stratonovich transformation to rewrite the quartic interaction term in a bilinear form, where the fermionic operators are now coupled to random auxiliary fields. Limiting oneself to purely superconducting order, the usual procedure for manipulating such a bilinear action is to separate the complex auxiliary field into its modulus and a pure phase part

before performing an expansion around the saddle point solution for the amplitude of this auxiliary field. This way of proceeding allows then to extract an effective action for the slow phase dynamics. This is distinctly different from the functional integral representation for the BFM which we have presented above and where from the very beginning and throughout such a procedure the fermionic pair fields are coupled to physically real bosonic modes which have their own proper dynamics. As we have seen in section III, the bosonic part of the action is treated within a coherent pseudo-spin representation where the dynamics of the pseudospin is parameterized by the time dependence of the spherical variables. The role of the dynamics within the spherical representation of the bosonic field and of the Berry phase term (which is a consequence of the quantum interference in the local pseudospin space) is a distinct feature of this BFM and presents specific differences with respect to the negative-U model. Such differences imply in particular that in the BFM case one has feedback effects between the amplitude and the phase fluctuations which are totally absent in an equivalent description of the negative-U Hubbard model. Such features are important because of intrinsic dissipation effects that eventually can change the nature of the transition and possibly are relevant for the emergence of a "bosonic metal" in proximity of the SIT, a feature which seems to be outside the framework of the negative-U Hubbard model.

Let us conclude this comparison of the BFM with similar models with a brief discussion on the Bose-Hubbard model and Josephson junction arrays. The main aspect which emerges as a common denominator in all those models, at least as far as the phenomenon of the SIT is concerned, is that the mechanism which is responsible for the degrading of the phase coherence is analogous and originates from the competition between the phase and charge degrees of freedom. In the Bose-Hubbard model this manifests itself as a competition between the boson hopping and their charge repulsion, while in the Josephson junction array scenario it appears as an interplay between the Josephson tunnelling amplitude and the charging energy.

In spite of this, at first sight, similarity between the BFM and those scenarios, we would like to stress that in a system where the dynamics is described by the BFM scenario, the interplay between the phase and amplitude fluctuations is intrinsically related to the coupling between the fermionic and bosonic degrees of freedom. This introduces amplitudes of the different processes at work in form of a non-trivial dependence on the microscopic parameters of the starting Hamiltonian. Furthermore, as we have seen in the above discussion of the effective action of the BFM, this model contains features which go beyond those which characterize the pure Josephson-type dynamics and which arise from the peculiar feedback between the fluctuations of the bosonic density (or amplitude pair field) and fluctuations of the phase. Last not least, the appearance of an intrinsic Berry phase-term in

the BFM scenario together with the effect of dissipation due to the fermionic dynamics, can give rise to an unconventional phenomenology when topological phases play a role, and especially in the presence of vortices.

VII. CONCLUSION

In this paper we examined the nature of a superconductor - insulator transition in a system of localized bosons and itinerant fermions coupled together via a pair exchange term. An effective action was derived from such a microscopic model, which, after integrating out the fermionic fields, could be phrased in terms of amplitude and phase fluctuations of the bosons. In order to make the presentation more familiar we discussed the action in a terminology frequently employed in connection with the study of Josephson junction arrays. We stress that our system does not necessarily imply any charged fermions and bosons.

Considering the phase-only part of the effective action it is fully equivalent to the quantum phase model for the Josephson junction arrays, discussed in terms of: (i) a *Josephson coupling* term, (ii) a *charging or capacitance* term and (iii) an *offset charge* term. Equivalent to that in our scenario is: (i) a boson hopping term, (ii) a term which takes into account the reduction in hopping amplitude due to a fluctuating local boson density arising from the intrinsic on-site exchange coupling between the bosons and the fermions and (iii) a chemical potential term controlling the bosonic concentration. New in the present study within already the lowest (phase-only) approach to the boson-fermion system is the intricacy of the dependence of the effective Josephson coupling, the capacitance term and the off-set charge term on the parameters of the initial Hamiltonian, i.e., the exchange coupling and the density of bosons (or else the total particle density). It turns out that within the lowest (phase-only) contribute a superconductor to insulator transition can not only be triggered by a change in the exchange coupling but also by a variation of the boson density. The latter presents evidently a certain interest from the experimental point of view and can possibly be tested in such transition occurring in optical lattices for ultracold fermi gases with Feshbach resonance pairing.

Apart from the phase-only part of the effective action we established the existence of an intrinsic Berry phase term, which arises from the hard core nature of the bosons and gives rise to an additional Magnus force when the system is such that topological ground state configurations like vortices are stabilized. This again is of potential interest for experiments since in principle one can change the sign of the Magnus force upon changing the density of bosons and thus deviating the motion of a vortex from one direction into the opposite one. These very preliminary results will be dealt with in greater detail in some future studies.

A further property of the Berry phase term is that it

gives rise to a bilinear coupling term between the phase and amplitude fluctuations, and is hence much more direct and relevant than similar terms in, for instance, scenarios based on the negative-U Hubbard model where they occur only at a much higher order in a corresponding loop expansion of the trace in the effective action. This again merits to be investigated in some detail with the aim to study the dissipation introduced by such amplitude phase coupling and its effect on the nature of the transition, in view of exploring the possibility of an intermediary bosonic metallic ground state.

Finally, we have dealt with a frequently asked question concerning the differences between the negative-U Hubbard scenario - mainly studied in connection with the BCS-BEC cross-over and the presently studied boson-fermion model. Far from being able to give a complete account for the major differences we found mainly two aspects which distinguish the physics of these models on a qualitative and robust level. The one is that in the negative-U Hubbard model a superconductor - insulator

transition can not take place at any finite coupling U , nor can such a transition be triggered by the change in particle concentration. The second point is that the ground states in the strong coupling limit of the two models are quite different: a highly degenerate ground state for the negative-U Hubbard model with excitations being controlled by an isotropic Heisenberg model when the hopping term is switched on. Contrary to that, for the boson fermion model the ground state is non-degenerate (corresponding to a ferro pseudo-magnetic Ising type system of singlets formed by bosons and fermion pairs), which, after switching on the fermion hopping term, gives rise to propagating domain like structures. The topological structures appearing in the ground state might have measurable consequences in the transport properties near the superconductor-insulator transition.

This and the other preliminary studies mentioned above will require a detailed analysis and will be discussed in future.

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