

# Numerical Renormalization Group at marginal spectral density: application to quantum tunneling in Luttinger liquids

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Many quantum mechanical problems (such as dissipative phase fluctuations in metallic and superconducting nanocircuits, or impurity scattering in Luttinger liquids) involve a continuum of bosonic modes with a marginal spectral density diverging as the inverse of energy. We construct a Numerical Renormalization Group in this singular case, with a manageable violation of scale separation at high energy, capturing reliably the low energy physics. The method is demonstrated by a non-perturbative solution over several energy decades for the dynamical conductance of a Luttinger liquid with a single static defect.

Bosonic description of fermionic systems, possibly subject to strong interactions, has a long history, ranging from phase fluctuations in superconducting circuits<sup>1,2</sup> following Josephson's initial ideas, to quantum transport in metallic grains<sup>3</sup> and in strongly correlated materials near the Mott transition,<sup>4</sup> where the phase conjugate to the electron charge is the relevant physical variable to understand the interplay of tunneling and Coulomb blockade. Another interesting example concerns one-dimensional electronic wires, the so-called Luttinger liquids (LL), where non-interacting plasmon modes provide a faithful representation of electronic density fluctuations.<sup>5-7</sup> Quite remarkably, all these different physical problems share very common features, because one can describe the electrons by a bosonic variable  $\Phi$  conjugate to the electronic charge transferred in a nanoscale junction (or also to the charge density in a wire), via a phase factor  $e^{i\Phi}$ . This mapping can be used for instance to represent the Josephson current in superconducting junctions (or also the fermionic fields in the bosonization language). This implies in turn that the nature of the phase dynamics determines the underlying physics: wild fluctuations of  $\Phi$  occur for instance in the presence of strong Coulomb blockade, leading to a rapid decay of the phase, and implying electronic localization.<sup>3,4</sup> In contrast, phase localization is associated to small fluctuations of the electron charge, characterizing dissipationless supercurrent<sup>1</sup> or Fermi liquid states.<sup>4</sup> The intermediate situation of soft (algebraic) phase decay leads to the well-known non-Fermi liquid features of a LL.<sup>5-7</sup>

In all these situations, great complexity arises due to the coupling of the bosonic mode to static disorder or dynamical defects, such as discrete Andreev levels<sup>8,9</sup> in superconducting weak links<sup>10</sup> or magnetic Kondo impurities in metallic junctions<sup>11,12</sup> and interacting unidimensional wires.<sup>13,14</sup> Focusing the discussion on the case of impurity effects in LL, but keeping this more general framework in mind, many technical and physical questions are still open to date, both in the original fermionic formulation and in the bosonic version of the problem. On the fermionic side, one needs to handle strong interactions within unidimensional wires together with the pres-

ence of exponentially small energy scales arising from the impurity,<sup>11,15</sup> for which powerful numerical methods have been developed in the past. The Density Matrix Renormalization Group<sup>16</sup> can tackle correlated wires, but on a linear energy scale only, which does not allow to really extract critical exponents;<sup>17</sup> the Numerical Renormalization Group (NRG)<sup>18,19</sup> can however deal with impurity physics on an exponential energy range, but only for uncorrelated Fermi liquids. A method that could incorporate both virtues would therefore be quite useful, which is the goal of this Letter.

Using the bosonic language, the description of interacting electrons by non-interacting bosons helps tremendously, but difficulties still arise. Apart from perturbative analysis (or fine tuning of the model parameters to allow an exact solution),<sup>5-7</sup> the analytical bosonization technique offers limited information on quantum impurity problems, because the physics crosses over from weak to strong coupling, for instance due to Kondo screening. Actually, all this complexity is already encoded by a static defect in LL, which drives the conductance from  $e^2/h$  to zero on an energy scale that can be exponentially small in the backscattering amplitude, a problem that has triggered substantial work, based on approximate analytical methods<sup>15,17,20-23</sup> or numerical techniques on a linear energy scale, such as quantum Monte Carlo.<sup>2,24</sup>

The idea we henceforth present here is to use recent developments of the bosonic NRG<sup>25-27</sup> in order to tackle numerically the phase fluctuation problem in a broad range of parameters, with possible extensions to dynamical defects. This however faces an immediate and seemingly intractable difficulty. Quite common to the Josephson effect in a dissipative environment,<sup>9</sup> to quantum tunneling in resistive circuits,<sup>12,28-30</sup> or to tunneling into LL,<sup>12,15,17,20-23</sup> is the marginal form of the bare local bosonic spectrum, given by the correlation function  $\mathcal{G}_{\Phi}^0(i\omega) = \frac{2\pi}{|\omega|}$  at imaginary frequency. The key step in the NRG procedure is the scale separation that results from a logarithmic discretization of the energy band  $\omega_n = \omega_c \Lambda^{-n}$ , with  $1 < \Lambda$ . A generalized power-law density of states (with exponent  $s \geq -1$  and high energy cutoff  $\omega_c$ ) of the form  $J(\omega) = 2\pi\omega_c^{-1-s}\omega^s\Theta(\omega_c - \omega)$  can be

considered both for fermionic<sup>31</sup> and bosonic models,<sup>25–27</sup> providing the following coupling strength of the states at energy  $\omega_n$ :

$$\gamma_n^2 = \int_{\omega_{n+1}}^{\omega_n} d\omega J(\omega) = 2\pi \frac{1 - \Lambda^{-(s+1)}}{s+1} \Lambda^{-n(s+1)}. \quad (1)$$

For all  $s > -1$ , the couplings  $\gamma_n^2$  decay exponentially with  $n$ , which allows an iterative diagonalization of the problem: the possibility of building progressively the Hilbert space from high to low energies is the reason behind the huge success of NRG to solve quantum impurity problems in a linear numerical effort.<sup>18,19</sup> However, the marginal case  $s = -1$  is special in the sense that the couplings  $\gamma_n^2 = 2\pi \log(\Lambda)$  do not decay anymore, invalidating clearly the whole scheme. We stress that we are considering quantum impurity Hamiltonians that depend explicitly on the phase factor  $e^{i\Phi}$ , and not on the spatial derivative of the bosonic mode,  $\partial_x \Phi$ . This latter case, which arises for instance in the so-called ohmic spin-boson model,<sup>25</sup> corresponds to the much simpler situation of linear spectrum ( $s = 1$ ) of the field  $\partial_x \Phi$ , and can be easily handled by the bosonic NRG.

This complete violation of scale separation for the marginal case  $s = -1$  seems to disqualify our proposed extension of the NRG. However, a free electron wire with constant density of states (described by the standard fermionic NRG at  $s = 0$ ) is equivalent to a free bosonic bath with  $s = -1$  due to the bosonization mapping, so that one may believe that the marginal situation could be tackled using some clever variant of the bosonic NRG. In order to move forward, let us investigate with greater detail the problem of tunneling in LL. The fermionic Hamiltonian reads in terms of second quantized left and right moving electron modes  $\psi_{L,R}^\dagger(x)$  at linear position  $x$  in the wire (omitting the electron spin for simplicity):

$$H = \int dx [iv_F \psi_L^\dagger \partial_x \psi_L - iv_F \psi_R^\dagger \partial_x \psi_R + g_2 \psi_R^\dagger \psi_R \psi_L^\dagger \psi_L - V_{bs} [\psi_R^\dagger \psi_L + \psi_L^\dagger \psi_R]_{x=0}] \quad (2)$$

where  $v_F$  is the Fermi velocity,  $g_2$  the short-range Coulomb repulsion between left and right moving electrons, and  $V_{bs}$  the impurity backward scattering amplitude at the  $x = 0$  location of the defect (forward impurity scattering and  $g_4$  interaction within a given Fermi point do not affect the physics, and were discarded.). The presence of the interaction term  $g_2$  clearly prevents a direct fermionic NRG solution of the model, which requires Fermi liquid leads. Yet, one can use the exact bosonization mapping<sup>5–7</sup> to re-express the electronic variables in terms of non-interacting collective charge density excitations  $\Phi(x)$  and conjugate field  $\Pi(x)$ . After standard manipulations<sup>5–7,15</sup> one obtains

$$H = \int \frac{dx}{8\pi} \{ [\Pi(x)]^2 + [\partial_x \Phi(x)]^2 \} - v \dagger \cos[\sqrt{K} \Phi(x=0)] \dagger \quad (3)$$

where normal ordering of the cosine operator, which will be crucial for the rigorous formulation of the NRG algorithm, has been emphasized. We have also introduced a small backscattering energy scale  $v \propto V_{bs}$  and the important Luttinger liquid parameter  $K = [(1 - g_2)/(1 + g_2)]^{1/2} \leq 1$ , into which all interaction effects have been encapsulated.

Let us now present how the bosonic NRG<sup>25</sup> can be tailored to address the impurity model (3), which has the form of a boundary Sine Gordon Hamiltonian. The derivation of the “star”-NRG follows the usual procedure<sup>19,25</sup> by considering the equivalent energy representation in terms of a continuum of canonical bosons  $a_\epsilon^\dagger$ :

$$H = \int_0^{\omega_c} d\epsilon \epsilon a_\epsilon^\dagger a_\epsilon - v \dagger \cos[\sqrt{K} \Phi] \dagger, \quad (4)$$

$$\Phi \equiv \Phi(x=0) = \sqrt{2} \int_0^{\omega_c} d\epsilon \frac{a_\epsilon^\dagger + a_\epsilon}{\sqrt{\epsilon}}. \quad (5)$$

The bosonic fields are then decomposed in Fourier modes ( $p \in \mathbb{Z}$ ,  $n \in \mathbb{N}$ ) on each interval  $\omega_{n+1} < \epsilon < \omega_n$  of width  $d_n = (1 - \Lambda^{-1})\Lambda^{-n}$ :

$$a_\epsilon^\dagger = \sum_{n,p} \frac{e^{i2\pi p\epsilon/d_n}}{\sqrt{d_n}} a_{n,p}^\dagger. \quad (6)$$

The first NRG approximation consists in neglecting all  $p \neq 0$  modes, keeping only the operators  $a_n^\dagger \equiv a_{n,0}^\dagger$  (this step becomes exact in the  $\Lambda \rightarrow 1$  limit<sup>19</sup>). This leads to the “star”-Hamiltonian:

$$H_S = \sum_{n=0}^{+\infty} \xi_n a_n^\dagger a_n - v \dagger \cos \left[ \sqrt{K} \sum_{n=0}^{+\infty} \frac{\gamma_n}{\sqrt{\pi}} (a_n^\dagger + a_n) \right] \dagger \quad (7)$$

with the “impurity” coupling strength already given in Eq. (1) by  $\gamma_n^2 = 2\pi \log(\Lambda)$  in the marginal case  $s = -1$ . The typical energy  $\xi_n$  in each shell is defined by:

$$\xi_n = \frac{1}{\gamma_n^2} \int_{\omega_{n+1}}^{\omega_n} d\omega \omega J(\omega) = \frac{1 - \Lambda^{-1}}{\log(\Lambda)} \omega_c \Lambda^{-n}. \quad (8)$$

As a benchmark of the discretization for the marginal case  $s = -1$ , one can easily compute from (7) the resulting approximation for the original Green’s function:

$$\mathcal{G}_{\Phi,\Lambda}^0(i\omega) = \frac{4}{1 - \Lambda^{-1}} \sum_{n=0}^{+\infty} \frac{\omega_c \Lambda^{-n}}{\omega^2 + \left[ \frac{1 - \Lambda^{-1}}{\log(\Lambda)} \right]^2 \omega_c^2 \Lambda^{-2n}} \quad (9)$$

which can be checked to converge exponentially fast at  $\omega \ll \omega_c$  to the exact result  $\mathcal{G}_{\Phi}^0(i\omega) = \frac{2\pi}{|\omega|}$  even for  $\Lambda = 2$  (we keep this standard value from now on). However, despite the clear exponential decay of the energies (8), the non-decreasing value of the couplings  $\gamma_n$  implies a violation of scale separation on *all* shells, and prevents the solution by iterative diagonalization of Hamiltonian (7).

The first key idea in successfully constructing the marginal bosonic NRG is to assume that the energy spectrum is also bounded from *below*:

$$J(\omega) = \frac{2\pi}{\omega} \Theta(\omega_c - \omega) \Theta(\omega - \omega_{\min}). \quad (10)$$

Clearly both the energies  $\epsilon_n$  and the couplings  $\gamma_n$  are not modified by this choice ( $\gamma_n$  still do not decay), and they are just cut off for  $n > n_{\min}$ , with  $\omega_{\min} = \omega_c \Lambda^{-n_{\min}}$ , so that nothing seems gained naively. We can however try to pursue with the second step of the standard NRG procedure, which amounts to the exact mapping on the Wilson chain.<sup>19,25</sup> This simple tridiagonalization procedure of Hamiltonian (7) leads to the following “chain” form in terms of new canonical bosons  $b_n^\dagger$ :

$$H_C = \sum_{n=0}^{+\infty} \left[ \epsilon_n b_n^\dagger b_n + t_n (b_n^\dagger b_{n+1} + b_{n+1}^\dagger b_n^\dagger) \right] - v \dagger \cos \left[ \sqrt{\frac{\eta_0 K}{\pi}} (b_0^\dagger + b_0) \right] \dagger \quad (11)$$

with the parameter  $\eta_0 = 2\pi \log(\omega_c/\omega_{\min})$ . Clearly, the impurity part of the chain Hamiltonian (11) breaks down for  $\omega_{\min} \rightarrow 0$ , owing to the divergence of  $\eta_0$ , but one can check numerically that the construction is valid for non-zero  $\omega_{\min}$ . The on-site energies  $\epsilon_n$  and hoppings  $t_n$  of the Wilson chain can indeed be obtained by numerical tridiagonalization of Eq. (7). For the value  $\omega_{\min} = 10^{-5}$  of the lower cutoff, these are plotted together with the star parameters on Fig. 1. The exponential decay of *both* chain

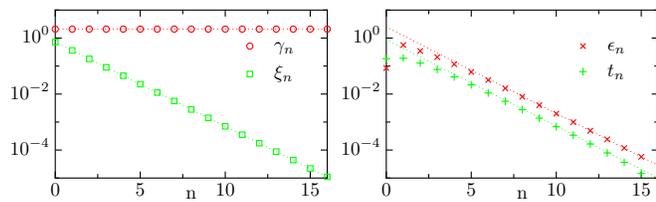


FIG. 1: (color online) Left panel: parameters  $\xi_n$  and  $\gamma_n$  of the star-NRG as a function of  $n$  for  $0 \leq n \leq n_{\min} = 16$ ; the coupling  $\gamma_n$  does not decay and violates scale separation on *all* shells. Right panel: parameters  $\epsilon_n$  and  $t_n$  of the chain-NRG; scale separation is only broken on the first shell, as seen by the initial increase of both parameters, before further exponential decay (shown by dotted lines as guides to the eye).

parameters  $\epsilon_n$  and  $t_n$  that we discover here is clearly a remarkable surprise, that enables the extension of the NRG to the marginal situation  $s = -1$ . This crucial feature comes at a small price, seen by the first increase of the chain parameters from site  $n = 0$  to site  $n = 1$ . Thus the maximal violation of scale separation in the star NRG presents a small remanence in the chain NRG, limited only to the first shell. Interestingly, the initial jump of the parameters is just proportional to  $\log(\omega_c/\omega_{\min})$ , so that the lower cutoff  $\omega_{\min}$  can be decreased on exponential scales without paying a huge numerical cost.

A last difficulty due to the unusual form of the impurity Hamiltonian (11) must be addressed. In the standard NRG,<sup>19,25</sup> only linear to quadrilinear operators are present in the Hamiltonian. However, the central physical role played by the phase factor  $e^{i\Phi}$  leads to a cosine term at the impurity site, hence to an operator of infinite order, which by the bosonization rules<sup>7</sup>

must also be normal ordered. For a generic operator  $\mathcal{O} = \dagger \cos[\alpha(b_0^\dagger + b_0)] \dagger$  this reads  $\mathcal{O} = \cos(\alpha b_0^\dagger) \cos(\alpha b_0) - \sin(\alpha b_0^\dagger) \sin(\alpha b_0)$ . Using the Fock states  $|m\rangle$  of the bosonic creation operator  $b_0^\dagger$  on the initial site  $n = 0$  of the Wilson chain, one obtains the matrix elements:

$$\langle m | \mathcal{O} | p \rangle = \sqrt{m!p!} \mathcal{R}e \sum_{k=0}^{\min(m,p)} \frac{(i\alpha)^{m+p-2k}}{(m-k)!(p-k)!k!}. \quad (12)$$

The construction of the impurity term in (11) proceeds by a truncation of the infinite Fock space on the initial Wilson site limited to states with occupation number less than a given  $N_0$ , and use of the matrix elements (12). Typically  $N_0 = 150$  ensures a good representation of the Hamiltonian. Each further site  $n > 0$  of the chain is described by a basis of  $N_b$  states<sup>25</sup> (we take  $N_b = 12$  here). At increasing  $n$ , the growing size of the total Hilbert space becomes rapidly unmanageable, and a truncation to  $N_{\text{trunc}}$  states is required (this approximation is common to all NRG schemes<sup>19</sup>). Typically  $N_{\text{trunc}} = 800$  was employed in all further computations, and we also set  $\omega_c = 1$  as the basic energy unit.

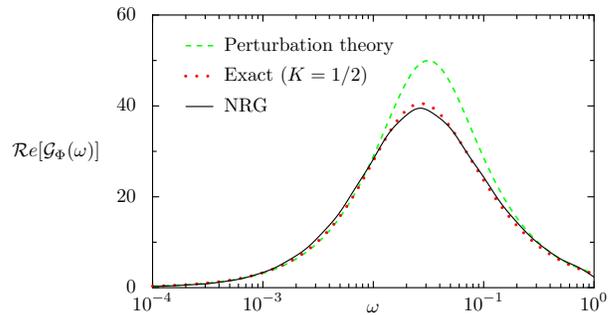


FIG. 2: (color online). Bosonic correlation function  $\text{Re}[\mathcal{G}_\Phi(\omega)]$  at real frequency  $\omega$  for the LL parameter  $K = 1/2$  comparing (bottom to top) the NRG to the exact result (13), and to the strong and weak interaction perturbation theory given respectively by (14) and (15) (these two expressions are by accident equivalent for  $K = 1/2$ , but nonetheless not exact).

In contrast to more complex extensions of impurity models with dynamical degrees of freedom (such as the Kondo model in a Luttinger liquid<sup>13,14</sup>), the present impurity problem benefits from several known limits, that allow to benchmark our numerics. For instance there exists an exact solution for the dynamical conductance<sup>5-7,15</sup> (in units of  $e^2/h$ ) at  $K = 1/2$ :

$$G^{\text{exact}}(\omega) = \frac{K\omega}{2\pi} \mathcal{R}e[\mathcal{G}_\Phi(\omega)] = \frac{1}{2} - \frac{\Omega}{2\omega} \text{atan} \left( \frac{\omega}{\Omega} \right) \quad (13)$$

where  $\Omega = e^\gamma v^2$ , with Euler’s constant  $\gamma$ , is the crossover energy at which the impurity cuts the chain (for  $K = 1/2$ ). Perturbation theory works also at strong interaction  $K \ll 1$ , in which case the self-consistent harmonic approximation applies:<sup>20</sup>

$$G^{\text{strong}}(\omega) = \frac{1}{2} \frac{\omega^2}{\omega^2 + (\Omega^*)^2} \quad (14)$$

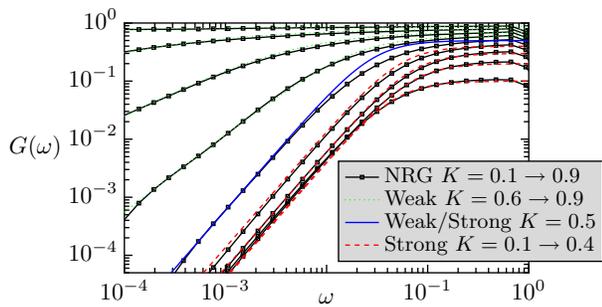


FIG. 3: (color online). Dynamical conductance  $G(\omega)$  in units of  $e^2/h$  for several values of the LL parameter  $K = 0.1, 0.2, 0.3, 0.4, 0.5, 0.6, 0.7, 0.8, 0.9$  obtained by the NRG (symbols, bottom to top). Comparison is made for  $K \leq 0.5$  to the strong interaction and for  $K \geq 0.5$  to the weak interaction limits. Perturbative methods cannot describe accurately the NRG solution beyond their applicability range.

with  $\Omega^* = 2\pi K v \frac{1}{1-K}$  the crossover scale. Finally, the limit of weak interaction  $1 - K \ll 1$  is also known from several approaches:<sup>5,6,15,17,21-23</sup>

$$G^{\text{weak}}(\omega) = \frac{K\omega^{\frac{2}{K}-2}}{\omega^{\frac{2}{K}-2} + (\Omega^*)^{\frac{2}{K}-2}}. \quad (15)$$

Fig. 2 compares our NRG data for  $K = 1/2$  with the ex-

act solution and perturbation theory, which shows the excellent convergence of the NRG and the sizeable discrepancies of both perturbative expansions. More systematic analysis for various  $K$  values in Fig. 3 demonstrates the progressive departure of the perturbative results from the numerical solution. The ability of the marginal bosonic NRG to describe non-perturbatively universal transport features with *high accuracy* should therefore make it a precious tool to test scaling behavior of impurity physics in LL.

To conclude, we have established an extension of the NRG to deal with the marginal situation of a density of states diverging as the inverse of energy. The potentially most promising applications of the NRG at marginal coupling concern the physics of dynamical impurities coupled to phase fluctuations, a large class of physical problems where no alternative analytical or numerical techniques exist to date. This development could allow to address many currently open issues, such as non-equilibrium transport with strong correlations (using a mapping onto equilibrium q-oscillator models<sup>32</sup>), Kondo physics in Luttinger liquids,<sup>12-14</sup> and ohmic dissipation in Andreev level qubits.<sup>8-10</sup>

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- <sup>1</sup> K. K. Likharev, *Dynamics of Josephson Junctions and Circuits* (Gordon, 1986).
  - <sup>2</sup> S. L. Lukyanov and P. Werner, *J. Stat. Mech.* **06**, 06002 (2007).
  - <sup>3</sup> G. Schön and A. D. Zaikin, *Phys. Rep.* **198**, 237 (1990).
  - <sup>4</sup> S. Florens and A. Georges, *Phys. Rev. B* **70**, 035114 (2004).
  - <sup>5</sup> T. Giamarchi, *Quantum physics in one dimension* (Oxford, 2003).
  - <sup>6</sup> A. O. Gogolin, A. A. Nersisyan, and A. M. Tsvelik, *Bosonization and strongly correlated systems* (Cambridge, 1998).
  - <sup>7</sup> J. von Delft and H. Schoeller, *Ann. Phys.* **7**, 225 (1998).
  - <sup>8</sup> A. Zazunov, V. S. Shumeiko, G. Wendin, and E. N. Bratus, *Phys. Rev. B* **71**, 214505 (2005).
  - <sup>9</sup> A. V. Shytov, arXiv:cond-mat/0001012.
  - <sup>10</sup> J.-D. Pillet *et al.*, *Nature* **6**, 965 (2010).
  - <sup>11</sup> A. C. Hewson, *The Kondo Problem to Heavy Fermions* (Cambridge, 1996).
  - <sup>12</sup> S. Florens, P. Simon, S. Andergassen, and D. Feinberg, *Phys. Rev. B* **75**, 155321 (2007).
  - <sup>13</sup> M. Fabrizio, and A. O. Gogolin, *Phys. Rev. B* **51**, 17827 (1995).
  - <sup>14</sup> R. Egger and A. Komnik, *Phys. Rev. B* **57**, 10620 (1998).
  - <sup>15</sup> C. L. Kane, and M. P. A. Fisher, *Phys. Rev. Lett.* **68**, 1220 (1992); *Phys. Rev. B* **46**, 7268 (1992).
  - <sup>16</sup> S. R. White, *Phys. Rev. Lett.* **69**, 2863 (1992).
  - <sup>17</sup> V. Meden *et al.*, *New J. Phys.* **10**, 045012 (2008).
  - <sup>18</sup> H.R. Krishna-murthy, J.W. Wilkins, and K.G. Wilson *Phys. Rev. B* **21**, 1003 (1980).
  - <sup>19</sup> R. Bulla, T. Costi and T. Pruschke, *Rev. Mod. Phys.* **80**, 395 (2008).
  - <sup>20</sup> F. Guinea, G. Gomez Santos, M. Sassetti, and M. Ueda, *Euro. Phys. Lett.* **30**, 561 (1995).
  - <sup>21</sup> D. Yue, L. I. Glazman, and K. A. Matveev, *Phys. Rev. B* **49**, 1966 (1994).
  - <sup>22</sup> T. Enss *et al.*, *Phys. Rev. B* **71**, 155401 (2005).
  - <sup>23</sup> D. N. Aristov, and P. Woelfle, *Euro. Phys. Lett.* **82**, 27001 (2008).
  - <sup>24</sup> K. Moon, H. Yi, C. L. Kane, S. M. Girvin, and M. P. A. Fisher, *Phys. Rev. Lett.* **71**, 4381 (1993).
  - <sup>25</sup> R. Bulla, N.-H. Tong, and M. Vojta, *Phys. Rev. Lett.* **91**, 170601 (2003); R. Bulla, H.-J. Lee, N.-H. Tong and M. Vojta, *Phys. Rev. B* **71**, 045122 (2005).
  - <sup>26</sup> M. T. Glossop, and K. Ingersent, *Phys. Rev. Lett.* **95**, 067202 (2005).
  - <sup>27</sup> A. Freyn and S. Florens, *Phys. Rev. B* **79**, 121102 (2009).
  - <sup>28</sup> I. Safi and H. Saleur, *Phys. Rev. Lett.* **93**, 126602 (2004).
  - <sup>29</sup> C. Mora, and K. Le Hur, *Nature Phys.* **6**, 697 (2010).
  - <sup>30</sup> Y. Hamamoto, T. Jonckheere, T. Kato, and T. Martin, *Phys. Rev. B* **81**, (2010) 153305.
  - <sup>31</sup> C. Gonzalez-Buxton, and K. Ingersent, *Phys. Rev. B* **57**, 14254 (1998).
  - <sup>32</sup> V. Bazhanov, S. Lukyanov, and A. Zamolodchikov, *Nucl. Phys. B* **549**, 529 (1999).