

## Interactions in Electronic Mach-Zehnder Interferometers with Copropagating Edge Channels

Luca Chirolli,<sup>1,2,\*</sup> Fabio Taddei,<sup>1</sup> Rosario Fazio,<sup>1</sup> and Vittorio Giovannetti<sup>1</sup>

<sup>1</sup>*NEST, Scuola Normale Superiore and Istituto Nanoscienze-CNR, I-56127 Pisa, Italy*

<sup>2</sup>*Instituto de Ciencia de Materiales de Madrid, CSIC, Sor Juana Ines de la Cruz 3, E-28049 Madrid, Spain*

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We study Coulomb interactions in the finite bias response of Mach-Zehnder interferometers, which exploit copropagating edge states in the integer quantum Hall effect. Here, interactions are particularly important since the coherent coupling of edge channels is due to a resonant mechanism that is spoiled by inelastic processes. We find that interactions yield a saturation, as a function of bias voltage, of the period-averaged interferometer current, which gives rise to unusual features, such as negative differential conductance, enhancement of the visibility of the current, and nonbounded or even diverging visibility of the differential conductance.

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**Introduction.**—Topological edge states in the integer quantum Hall effect [1] represent an ideal playground for testing the coherence of electronic systems at a fundamental level. The harnessing of edge states as quasi-one-dimensional (1D) chiral electronic waveguides has allowed the successful realization of a number of electronic interferometric setups, such as those of Mach-Zehnder [2–5] and Fabry-Perot [6]. These devices exploit counterpropagating edge states localized at opposite sides of a Hall bar, which are brought in contact and mixed by quantum point constrictions (QPCs) that mimic the effect of optical beam splitters (BSs). In Mach-Zehnder interferometer (MZI) setups, the chirality of electron propagation makes necessary the adoption of nonsimply connected, Corbino-like geometries, which limit the flexibility of the devices. In these systems, electron-electron ( $e$ - $e$ ) interactions are in general responsible for dephasing via the coupling with external edge channels [7,8], which manifests as a reduction of the visibility of the Aharonov-Bohm (AB) oscillations as a function of the bias voltage. In particular, puzzling behaviors have been reported [3–5] in the finite bias response of MZIs, in which the visibility presents a lobelike structure and phase rigidity [7,9–15].

An alternative MZI scheme inducing coherent mixing between edge states copropagating at the *same* boundary of the sample has been suggested [16] as a more flexible architecture, which allows multiple device concatenation [16,17] and represents an ideal candidate for implementation of dual-rail quantum computation schemes [18,19]. As schematized in Fig. 1, in such a setup the BS transformations are now implemented through series of top gates, organized in arrays of periodicity  $\lambda$  tuned to compensate for the difference between the momenta of the copropagating channels (the inner channel  $i$  and the outer channel  $o$  of the figure); see Ref. [20], where the first experimental realization of such BSs is reported, and Ref. [21]. The corresponding AB phase difference is introduced instead by separating the edge modes in the region between the

two BSs through the action of a central top gate [16]. A similar interferometer, but featuring no modulation, has been realized in Ref. [22]. The effects of  $e$ - $e$  interactions in these MZIs are likely to be qualitatively different from those observed in the Corbino-like settings of Refs. [2–5]. Indeed, whereas in the latter the direct coupling between the channels that enter the interferometer can be neglected, in the regions where they are coherently mixed (the QPCs), in the scheme of Fig. 1, this is no longer possible due to the non-negligible spatial extension of the top gate arrays. This implies a strong interplay between coherent mixing and interactions that might impair the MZI response. The aim of the present work is to target such interplay.

As detailed in the following, we show that the interedge current  $I$  measured at the output of the setup of Fig. 1 possesses a strong nonlinear dependence in the bias voltage  $V$  that, while still exhibiting AB oscillations, leads to saturation of the associated mean value averaged over the AB phases. Hence, the device presents negative differential conductance with unbounded visibility. Such anomalous behavior occurs since interactions enable inelastic scattering that spoils, for increasing voltages, the coherence

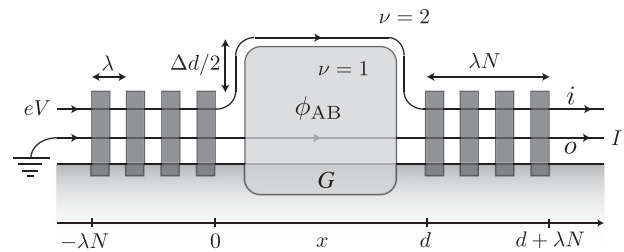


FIG. 1. Sketch of the MZI. Two sets of  $N$  top gates arranged in arrays with periodicity  $\lambda$  and separated by the distance  $d$  represent the  $L$  and  $R$  BSs of the interferometer. A central top gate  $G$  locally lowers the filling factor to  $\nu = 1$  and separates the two edge states ( $i$  and  $o$ ), which experience a path length difference  $\Delta d$  and acquire an AB phase  $\phi_{AB}$ .

needed for interedge coupling to occur at the BSs. Furthermore, as long as the interactions between the edge channels can be neglected in the region between the two BSs, we also observe that for large voltages the visibility of current gets enhanced with respect to the noninteracting case.

*Model.*—Before discussing the role of  $e$ - $e$  interactions in the device of Fig. 1, we find it useful to briefly review the basic properties of the scheme in the noninteracting case. The underlying idea [20] is to implement BS transformations between two copropagating channels ( $i$  and  $o$ ), via the action of a pair of arrays of top gates (see Fig. 1), which are spatially modulated at periodicity  $\lambda$ . Following Refs. [20,21], we describe them through potentials of the form  $t_L(x) = \bar{t}_L \sin^2(\pi x/\lambda)$ , for  $-\lambda N < x < 0$ , and  $t_R(x) = \bar{t}_R \sin^2[\pi(x-d)/\lambda]$ , for  $d < x < d + \lambda N$  ( $N$  being the number of elements of a single array whereas  $d$  is the distance between the two sets as measured according to the coordinates of the channel  $o$ ). With introduction of the difference  $\Delta k = k_i - k_o$  between the momenta  $k_i$  and  $k_o$  of the two edges, the tunneling amplitude at a given BS can then be expressed as  $\bar{t}_\alpha \text{sinc}[\lambda N(\Delta k - 2\pi/\lambda)/2]$ , at lowest orders in  $\bar{t}_\alpha$  (here  $\alpha = L, R$  indicates the left and right BS, respectively) [20]. In this scenario, Mach-Zehnder interferences can be observed by introducing between the two BSs a top gate that locally lowers the filling factor to  $\nu = 1$  and diverts the inner edge state toward the interior of the mesa [16]. This way, the two channels are guided along paths of difference lengths,  $d_o = d$  and  $d_i = d + \Delta d$ , thus acquiring an AB phase difference  $\phi_{AB}$  proportional to the magnetic field  $B$  and to the area enclosed by the path and a dynamical phase. The transmission probability  $T(\epsilon)$  at energy  $\epsilon$  of the MZI, from the inner channel on the left to outer channel on the right, is then given by  $T(\epsilon) \propto [|\bar{t}_L|^2 + |\bar{t}_R|^2 + 2|\bar{t}_L \bar{t}_R| \cos(\phi_{AB} + \epsilon \Delta d / v_F)] S^2$ , where  $v_F$  is the group velocity of edge states and  $S = \text{sinc}[\lambda N(\Delta k - 2\pi/\lambda)/2]$  is the BS amplitude, which is optimal when the resonant condition  $\Delta k = 2\pi/\lambda$  is met. A bias voltage  $V$  applied between channels  $i$  and  $o$  gives rise to a zero-temperature current  $I(V) \propto \int_0^{eV} d\epsilon T(\epsilon)$ , whose associated visibility  $\mathcal{V}_I = (\max_\phi I - \min_\phi I) / (\max_\phi I + \min_\phi I)$  of the AB oscillations amounts to  $\mathcal{V}_I = \mathcal{V}_\sigma^{(0)} |\text{sinc}(eV \Delta d / 2v_F)|$ , with  $\mathcal{V}_\sigma^{(0)} = 2|\bar{t}_L \bar{t}_R| / (|\bar{t}_L|^2 + |\bar{t}_R|^2)$  being the oscillation visibility of the differential conductance  $\sigma = dI/dV$ .

To analyze the effect of interactions in such a setup, we describe the linearly dispersing electronic excitations around the Fermi energy by means of two chiral fermion fields  $e^{ik_m x} \psi_m(x)$ , with  $m = i, o$ , each propagating at mean momentum  $k_m$ . The kinetic Hamiltonian can then be written as ( $\hbar = 1$ )  $H_{\text{kin}} = -iv_F \sum_m \int dx \psi_m^\dagger \partial_x \psi_m$ . The action of the BSs are instead assigned by means of the tunneling Hamiltonian  $H_{\text{tun}} = \sum_{\alpha=L,R} (A_\alpha + A_\alpha^\dagger)$ , where  $A_\alpha$  describe the action of the gate arrays potentials and are defined as  $A_L = \int dx t_L(x) e^{i\Delta k x} \psi_o^\dagger(x) \psi_i(x)$ ,  $A_R = \int dx t_R(x) e^{i\Delta k x} \psi_o^\dagger(x) \psi_i(x + \Delta d)$  with  $t_L(x)$ ,  $t_R(x)$ , and

$\Delta d$  introduced previously. In these expressions, the local phase shift  $e^{i\Delta k x}$  accounts for the resonant behavior of the MZI. The AB phase of the setup (together with a dynamical phase term  $k_i d_i - k_o d_o$ ) is instead included in the tunneling amplitude of the right beam splitter, i.e.,  $\bar{t}_R \rightarrow e^{i\phi_{AB}} \bar{t}_R$ .

As far as  $e$ - $e$  interactions are concerned, an electron propagating in one edge channel can interact with all the electrons in the Fermi seas of both channels. Although the precise form of the interaction potential is unknown, screening provided by top gates makes it reasonable to assume a short-range density-density interaction. The latter, however, need not be uniform in the whole sample: as a matter of fact, whereas intrachannel couplings are present everywhere, the interchannel interactions depend on the edge channel spatial separation, which in our setup varies strongly from place to place (see Fig. 1). In particular, in the region between the BSs edge states are brought far apart by the central top gate  $G$ , and it is reasonable to assume the interchannel coupling to be off. Interchannel interactions, however, cannot be excluded in the regions before and after  $G$ , where the tunneling term  $H_{\text{tun}}$  is active. Indicating with  $\rho_m = \psi_m^\dagger \psi_m$  the 1D density operator in channel  $m = i, o$ , we account for these effects by introducing an interchannel  $e$ - $e$  coupling  $H_{\text{inter}} = \int dx \int dx' : \rho_o(x) U(x, x') \rho_i(x') :$  characterized by a coordinate dependent potential  $U(x, x')$ , which nullifies in the central top gate region (i.e.,  $0 \leq x \leq d$ ) and which approaches the short-range behaviors  $2\pi g \delta(x - x')$  and  $2\pi g \delta(x - x' - \Delta d)$  on the lhs part (i.e.,  $x \leq 0$ ) and on the rhs part ( $x \geq d$ ) of the setup, respectively (the transition between these regions being smooth). In these expressions, the interaction strength  $g$  has the dimension of a velocity whereas, similarly to  $H_{\text{tun}}$ , the parameter  $\Delta d$  accounts for the relative coordinate shift experienced by the inner edge with respect to the outer. A short-range intrachannel coupling term of the form  $H_{\text{intra}} = \pi u \sum_m \int dx : \rho_m^2(x) :$  is also considered (in this case, however, no coordinate dependence is assumed).

*Methods.*—Including the  $e$ - $e$  interaction, the response of the interferometer driven out from equilibrium by a bias voltage  $\mu_i - \mu_o = eV$  applied between the two edge states will be analyzed at lowest order in the tunneling Hamiltonian  $H_{\text{tun}}$  (an exact analytical treatment being impossible). At second order in the amplitudes  $\bar{t}_\alpha$ , this allows us to express the current flowing through the outer edge as  $I(V) = e \sum_{\alpha,\beta} \int dt \langle [A_\alpha^\dagger(t), A_\beta(0)] \rangle$ , where  $A_\alpha(t) = e^{iH_0 t} A_\alpha e^{-iH_0 t}$  is the tunneling term evolved through the kinetic and interaction components of the system Hamiltonian, i.e.,  $H_0 = H_{\text{kin}} + H_{\text{intra}} + H_{\text{inter}}$ , while expectation values are taken with respect to the ground state of the system biased by the chemical potential difference  $eV$  [7,23].

Expanding the summation over  $\alpha$  and  $\beta$ , we recognize the presence of three contributions:  $I = I_L + I_R + I_\phi$  with  $I_\alpha = e \int dt \langle [A_\alpha^\dagger(t), A_\alpha(0)] \rangle$ , for  $\alpha = L, R$  being direct

terms which do not depend upon the relative phase accumulated by the electrons when traveling through the MZI, and with  $I_\Phi = e \int dt \langle [A_L^\dagger(t), A_R(0)] \rangle + \text{c.c.}$  being a cross term, which is sensitive to the  $AB$  phase. Explicit expressions for these quantities are obtained by means of the two-point electron and hole Green's functions  $\mathcal{G}_m^e(x, t; x') = \langle \psi_m(x, t) \psi_m^\dagger(x', 0) \rangle$  and  $\mathcal{G}_m^h(x, t; x') = \langle \psi_m^\dagger(x, t) \psi_m(x', 0) \rangle$ , which we calculate by bosonization of the Hamiltonian  $H_0$ . Following the formalism of Refs. [7,12,24], we introduce chiral bosonic fields  $\phi_m$ , which satisfy the commutation rules  $[\phi_m(x), \phi_{m'}(x')] = i\pi\delta_{m,m'} \text{sgn}(x-x')$  [24,25] and express the fermion fields as  $\psi_m = (\hat{F}_m/\sqrt{2\pi a})e^{2\pi i\hat{N}_m x/L}e^{-i\phi_m}$  where  $a$  is a cutoff length that regularizes the theory at short wavelengths,  $\hat{F}_m$  are the Klein operators,  $\hat{N}_m = \int dx: \rho_m(x):$  are the total number operators of the edge states, and finally  $L$  is the edge quantization length. Observing that  $\rho_m = (1/2\pi)\partial_x\phi_m + \hat{N}_m/L$ , the kinetic and intrachannel Hamiltonian (apart for an irrelevant term proportional to  $\sum_m \hat{N}_m$ ) becomes  $H_{\text{kin}} + H_{\text{int}}^{\text{tra}} = (v/4\pi)\sum_m \int dx (\partial_x\phi_m)^2 + H_C^{\text{tra}}$  where  $v = v_F + u$  is the renormalized edge group velocity [7,12,24] and where  $H_C^{\text{tra}} = \pi v \sum_m \hat{N}_m^2/L$  is a capacitive contribution. Vice versa, by exploitation of the smooth variation assumption of the potential  $U(x, x')$  and the fact that  $L$  is the largest length in the problem, the interchannel interaction term yields

$$H_{\text{int}}^{\text{inter}} \simeq g \int_{-\infty}^0 \frac{dx}{2\pi} (\partial_x\phi_o)(\partial_x\phi_i) + g \int_d^\infty \frac{dx}{2\pi} (\partial_x\phi_o(x))(\partial_x\phi_i(x+\Delta d)) + H_C^{\text{inter}}, \quad (1)$$

with  $H_C^{\text{inter}} = (2\pi g/L)\hat{N}_o\hat{N}_i$  being a cross capacitive contribution. In contrast with Ref. [12], the interchannel interaction Eq. (1) is specifically active only in the tunneling regions. This will lead to qualitatively different results. The Hamiltonian  $H_0$  can now be brought to a diagonal form by solving the eigenvector equation  $[H_0, \hat{\gamma}_\pm(\epsilon)] + \epsilon\hat{\gamma}_\pm(\epsilon) = 0$ , which defines bosonic energy eigenmodes  $\hat{\gamma}_\pm(\epsilon)$ . Accordingly, we obtain

$$\phi_m(x, t) = \int_0^\infty \frac{d\epsilon}{\sqrt{\epsilon}} e^{-i\epsilon(t-i\tau)} \sum_{s=\pm} \varphi_m^s(x, \epsilon) \hat{\gamma}_s(\epsilon) + \text{h.c.},$$

where the wave functions  $\varphi_m^s$  satisfy the relation  $\sum_s \int_0^\infty (d\epsilon/\epsilon) \text{Im}[\varphi_m^s(x, \epsilon)\varphi_{m'}^s(x', \epsilon)^*] = (\pi/2)\delta_{m,m'} \text{sgn}(x-x')$ , and  $1/\tau > 0$  is an energy cutoff related to  $a$  via the noninteracting dispersion  $\epsilon = v_F q$ , i.e.,  $a/\tau = v_F$ . Solving the equations of motion by requiring only continuity of the wave functions, we find  $\varphi_m^\pm = f_\pm(x)$  for  $x < 0$ ,  $\varphi_m^\pm = C_\pm e^{i\epsilon x/v}$  for  $0 < x < d_m$ , with  $C_+ = 1$ ,  $C_- = 0$ , and  $\varphi_m^\pm = e^{i\epsilon d_m/v} f_\pm(x - d_m)$  for  $x > d_m$ , in terms of the symmetric and antisymmetric combinations

$f_\pm(x) = (e^{i\epsilon x/v_+} \pm e^{i\epsilon x/v_-})/2$ . Two new velocities enter the problem, a fast charged mode that propagates at  $v_+ = v + g$  and a slow neutral mode that propagates at  $v_- = v - g$ . With exploitation of these results, the electron and hole Green's functions can be finally written as  $\mathcal{G}_m^e = e^{-i\mu_m[t-(x-x')/v_-]} \mathcal{G}_m(x, t; x')$  and  $\mathcal{G}_m^h = e^{i\mu_m[t-(x-x')/v_-]} \mathcal{G}_m(x, t; x')$ , with the zero-bias Green's function  $\mathcal{G}_m(x, t; x') = (1/2\pi a) \langle e^{i\phi_m(x,t)} e^{-i\phi_m(x',0)} \rangle$ , which because of the quadratic nature of the bosonized Hamiltonian  $H_0$  can easily be computed in terms of the wave functions  $\varphi_m^s$ . In particular, to compute the direct current terms  $I_L$  and  $I_R$  we only need the correlation functions in the BS regions:  $x, x' < 0$  for the left BS and  $x, x' > d_m$  for the right BS. At zero temperature, for these combinations we find  $\mathcal{G}_m = (i/2\pi v_F) X_+^{-1/2} X_-^{-1/2}$ , with  $X_s = (x-x')/v_s - t + i\tau$ , in agreement with Ref. [12] (for finite temperatures see Ref. [26]). The term  $I_\Phi$  is obtained instead through crossed combinations. In particular for  $x > d_m$  and  $x' < 0$  we get  $\mathcal{G}_m = (i/2\pi v_F) \prod_{s=\pm} [X_s + d_m(1/v - 1/v_s)]^{-1/2}$  and analogously for  $x' > d_m$  and  $x < 0$ , with the replacement  $d_m \rightarrow -d_m$ .

*Currents.*—With the help of the Green's functions, the  $I$ - $V$  characteristics of the setup can now be explicitly computed. In particular, for the direct contributions of the current one gets

$$I_\alpha(V) = \frac{e}{2\pi} n_F^2 |\bar{t}_\alpha|^2 \int_0^{eV} d\epsilon S^2(\epsilon/\epsilon_c), \quad (2)$$

with the resonance function  $S(x) = \text{sinc}[\lambda N(\Delta k - 2\pi/\lambda)/2 - x/2]$  and the density of states at the Fermi energy  $n_F = \lambda N/(2v_F)$ . The energy scale  $\epsilon_c = [\lambda N(1/v_- - 1/v_+)]^{-1}$  is associated to the difference in time of flight for propagation of the bosonic excitations at speeds  $v_+$  and  $v_-$  through the BS of length  $\lambda N$ : in the absence of interactions it diverges. We notice that whereas for a low bias  $I_\alpha$  is linear in  $V$ , for a bias larger than  $\epsilon_c$  it saturates to the constant value  $I_\alpha^{\text{asy}} = \pi G_0 n_F^2 |\bar{t}_\alpha|^2 \epsilon_c$ . This is shown by the dashed lines in Fig. 2, top panel, where  $I_\alpha$  is plotted as a function of  $V$  for different values of  $g$ . Such behavior is due to the fact that inelastic processes, which are induced by the interaction and increase with increasing voltage bias, spoil the resonant coherent effect that is responsible for the transfer of charges between the two edge channels in the BS, thus suppressing  $I_\alpha$ .

The cross term contribution  $I_\Phi$  to the current is obtained instead through integration over the branch cuts [12] of the Green's functions in the cross region. Introducing

$$\mathcal{A} = \int_0^{eV} d\epsilon e^{i\varphi_V(\epsilon)} S^2(\epsilon/\epsilon_c) J_0\left(\frac{\epsilon}{2\epsilon_{\text{dyn}}}\right) J_0\left(\frac{eV - \epsilon}{2\epsilon_{\text{dyn}}}\right),$$

together with the phase  $\varphi_V(\epsilon) = \epsilon/\epsilon_c + (\Delta k - 2\pi/\lambda)\lambda N + eV[\Delta d/(2v_-) + (d + \Delta d/2)(1/v - 1/v_-)]$ , we obtain

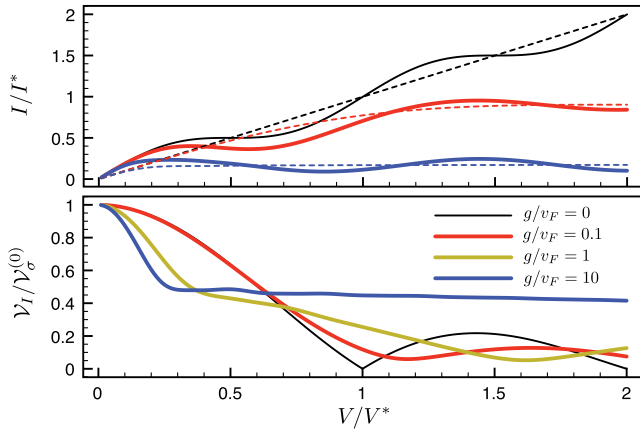


FIG. 2 (color online). Response of the MZI versus bias  $V$  [in units of  $V^* = 2\pi v_F/(e\Delta d)$ ] for different values of the interaction parameter  $g/v_F$ . We have assumed  $u = g$ ,  $\phi_{AB} = 0$ ,  $|\bar{t}_L| = |\bar{t}_R| = |\bar{t}|$ , and that the resonant condition  $\Delta k = 2\pi/\lambda$  is satisfied. Top: total current  $I$  (full curves) and BS current (dashed curves) in units of  $I^* = n_F^2 |\bar{t}|^2 e^2 V^*/\pi$ . Bottom: visibility of the AB oscillations in the current  $\mathcal{V}_I$  in units of  $\mathcal{V}_\sigma^{(0)}$ . According to Ref. [20], we set  $\lambda N/\Delta d = 3$  and  $d/\Delta d = 1$ .

$$I_\Phi(V) = \frac{e}{2\pi} n_F^2 2|\bar{t}_L \bar{t}_R| |\mathcal{A}| \cos[\phi_{AB} + \arg(\mathcal{A})],$$

where  $J_0[x]$  is the Bessel function of the first kind and  $\epsilon_{\text{dyn}} = v/\Delta d$ . The latter is a new energy scale associated to the dynamical phase difference acquired by the electrons in the interference region between the BSs, where interchannel interaction is absent and excitations move at speed  $v$ . Due to the bias dependence of  $\arg(\mathcal{A})$ , the current  $I_\Phi$  shows oscillations versus bias around zero.

The overall  $I$ - $V$  characteristics (full lines plotted for different values of  $g$  in Fig. 2, top panel) show an oscillating behavior that becomes nonmonotonic for large values of the bias  $V$ . The associated visibility is  $\mathcal{V}_I(V) = \mathcal{V}_\sigma^{(0)} |\mathcal{A}| / \int_0^V d\epsilon S^2(\epsilon/\epsilon_c)$ , where  $\mathcal{V}_\sigma^{(0)}$ , as in the absence of interactions, provides the upper bound ( $\mathcal{V}_I \leq \mathcal{V}_\sigma^{(0)} \leq 1$ ). Plots of  $\mathcal{V}_I(V)$  for different values of  $g$  are shown in the bottom panel of Fig. 2. For  $g/v_F = 0.1$  (red curve), weakly interacting regime,  $\mathcal{V}_I(V)$  closely follows the noninteracting case (black curve) for  $V \leq V^* = 2\pi v_F/(e\Delta d)$  and thereafter oscillating without reaching zero. In the strongly interacting regime  $g \gg v_F$ , the scale  $\epsilon_c$  already dominates at relatively low biases, leading to a completely different behavior. As shown by the blue curve ( $g/v_F = 10$ ),  $\mathcal{V}_I(V)$  decreases with the bias very rapidly up to  $eV \leq \epsilon_c$  and very slowly thereafter, on the scale  $\epsilon_{\text{dyn}} [2\pi\epsilon_c/(eV^*)] \approx 0.32$  and  $2\pi\epsilon_{\text{dyn}}/(eV^*) \approx 11$ , with the parameters used in Fig. 2]. The almost constant value exhibited by the curve is the consequence of the fact that the current oscillates around a constant value, with an amplitude of oscillations that decreases very slowly with bias. Interestingly, in this regime the visibility in the interacting

case is higher than that in the noninteracting case. Also very peculiar is the case of a symmetric interferometer  $\Delta d = 0$ , where  $\epsilon_{\text{dyn}}$  diverges and the visibility becomes  $\mathcal{V}_I^{\Delta d=0} = \mathcal{V}_\sigma^{(0)} |\mathcal{F}_2(V)|/\mathcal{F}_0(V)$ , where  $\mathcal{F}_p(V) = \int_0^{eV/(2\epsilon_c)} dx e^{ixp} \text{sinc}^2(x)$ . We find that the ratio  $\mathcal{V}_I^{\Delta d=0}/\mathcal{V}_\sigma^{(0)}$  goes, for  $eV \gg \epsilon_c$ , asymptotically to the finite value  $2\ln(2)/\pi$ . Therefore, a symmetric device overperforms with respect to an asymmetric one for large voltages.

The highly nonmonotonic behavior of  $I(V)$  detailed above is best understood by studying the differential conductance  $dI/dV$ . In particular, its direct contribution amounts to  $dI_\alpha/dV = (e/2\pi)n_F^2 |\bar{t}_\alpha|^2 S^2(eV/\epsilon_c)$ , and it shows how  $e$ - $e$  interactions effectively give rise to an interaction- and bias-dependent shift of the resonance condition

$$\Delta k = \frac{2\pi}{\lambda} + \frac{2g}{v_+ v_-} eV. \quad (3)$$

As a result,  $dI_\alpha/dV$  becomes negligible beyond an applied voltage on order of  $\epsilon_c$ , which explains the saturation of the period-averaged current in Fig. 2 in the strongly interacting regime. At the same time, a non-negligible cross term  $dI_\Phi/dV$  results in an overall behavior of the total differential conductance  $dI/dV$ , which exhibits regions of negative values. Furthermore, while the associated visibility of the AB oscillations in the presence of interactions is known to exhibit values greater than 1 [12], for our setup this quantity diverges for values of bias  $V$  such that the direct contributions to the differential conductance vanish. The above peculiar behavior can be harnessed to improve the performances of the MZI. For example, via Eq. (3), the bias  $V$  can be used as a knob for fine tuning of the resonance condition, as the precise value of the momentum difference  $\Delta k$  is *a priori* unknown. We finally note that for systems that do not feature a modulation of the tunneling ( $\lambda \rightarrow \infty$ ), Eq. (3) predicts that charge transfer can be achieved for sufficiently high bias, as an interaction-driven resonance is met for a threshold voltage  $eV_{\text{th}} = v_F \Delta k$ . This picture shares analogies with the experimental findings of Deviatov *et al.* [22], who reported interedge transport and AB oscillations only beyond a threshold bias.

*Conclusions.*—We have shown that  $e$ - $e$  interactions reduce the performances of a MZI with copropagating edge states. There are, however, striking differences with respect to nonsimply connected MZI architectures. Indeed, while in the latter case the interferometer edge channels are coupled to additional modes that carry information away from the system [3–5, 12], in our setup they only interact among themselves and the information is redistributed in the system. The major impact is the spoiling of the resonant tunneling condition that realizes the BS (the oscillating component of the interferometer current being only marginally affected because interedge interactions are negligible in the interference region). This leads to the



unexpected result that strong interactions yield a reduction of the current visibility for small voltages, but an enhancement for larger voltages, with respect to the noninteracting case. Furthermore, the differential conductance can become negative in some voltage ranges, whereas its visibility can take large values or even diverge.

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\*luca.chirolli@icmm.csic.es

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- [26] At finite temperature  $T$ , the Green's function in the BS regions is modified to give  $\mathcal{G}_m = (i/2\pi v_F) \prod_{s=\pm} [\pi T / \sinh(\pi T X_s)]^{1/2}$  and analogously in the cross regions. In the resulting expression for the total current, the integration over the bias window is regulated by Fermi functions, which smooth the entire energy and flux dependent integrand [in particular the resonance function  $S(x)$ ] around 0 and  $eV$ , on the scale  $T$ . In the amplitude  $\mathcal{A}$  of the interference term, the Bessel functions are replaced by hypergeometric functions, which cause a further reduction of the visibility over an energy scale  $\epsilon_{\text{dyn}}$ . As a consequence, as in Refs. [5,12], finite-bias and finite-temperature dephasing are determined by the same energy scales, namely,  $\epsilon_c$  and  $\epsilon_{\text{dyn}}$ .