# Fractionally Charged Zero-Energy Single-Particle Excitations in a Driven Fermi Sea

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A voltage pulse of a Lorentzian shape carrying half of the flux quantum excites out of a zero-temperature Fermi sea an electron in a mixed state, which looks like a quasiparticle with an effectively fractional charge $e/2$. A prominent feature of such an excitation is a narrow peak in the energy distribution function lying exactly at the Fermi energy $\mu$. Another spectacular feature is that the distribution function has symmetric tails around $\mu$, which results in a zero-energy excitation. This sounds improbable since at zero temperature all available states below $\mu$ are fully occupied. The resolution lies in the fact that such a voltage pulse also excites electron-hole pairs, which free some space below $\mu$ and thus allow a zero-energy quasiparticle to exist. I discuss also how to address separately electron-hole pairs and a fractionally charged zero-energy excitation in an experiment.

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*Introduction.*—The recent realization of a triggered single-electron source [1–10] opens a new era for coherent electronics [11–18] by allowing it to become quantum much like quantum optics. The analogues of the famous quantum optics effects were successfully demonstrated with single electrons in solid-state circuits such as the partitioning of electrons [7,19–21] in Hanbury Brown–Twiss geometry and quantum-statistical repulsion of electrons [7,22] in Hong-Ou-Mandel geometry. Tomography of a single-electron state [23] and a preparation of few-electron Fock states [20,24,25] are already reported.

An essential difference from quantum optics is that single electrons are injected into electron waveguides, such as, for instance, quantum Hall edge channels [26–28], which contain other electrons [29]. During such an injection the source can excite an electron system and the resulting excitations can mask injected electrons. However, if the protocol of injection is properly chosen [30], no spurious excitations appear. This was clearly demonstrated theoretically [31–33] and experimentally [7] in the case where single electrons are excited by applying a voltage pulse $V(t)$ across a ballistic conductor. It was shown that a voltage pulse of a Lorentzian shape with a quantized flux, $\varphi \equiv (e/\hbar)\int dtV(t) = 2\pi n$ (where $e$ is the electron charge, $\hbar$ is Planck's constant, and $n$ is an integer), excites only $n$ electrons (or holes, if $n < 0$) with no accompanying electron-hole pairs. These excitations were named levitons [7]. If the flux is not quantized, $\varphi \neq 2\pi n$, then what is excited is rather a messy state with a divergent number of quasiparticles, both electrons and holes.

Here, I show, however, that the flux $\varphi = \pi$ is special. The Fermi sea excited by a Lorentzian voltage pulse with a half-integer flux hosts an exotic single-particle excitation which cannot exist in equilibrium; see Fig. 1. Such an excitation has an effective charge $e/2$; hence, I call it a *half-leviton* (HL). Importantly, an electron-hole state (which is also excited because the flux is not quantized) is indispensable to the existence of HLs. This is so because the state of a half-leviton is a superposition of states with energies lying on both sides of the Fermi energy $\mu$, below and above it. At zero temperature, all of the states below the Fermi energy are fully occupied and only excited holes (belonging to electron-hole pairs) free some states below $\mu$ and allow a half-leviton to be formed.

Importantly, a half-leviton has zero energy. This allows it to annihilate effectively (without breaking a phase coherence) its antiparticle, which is excited by a voltage pulse carrying a flux of the opposite sign; see Fig. 2. Such a

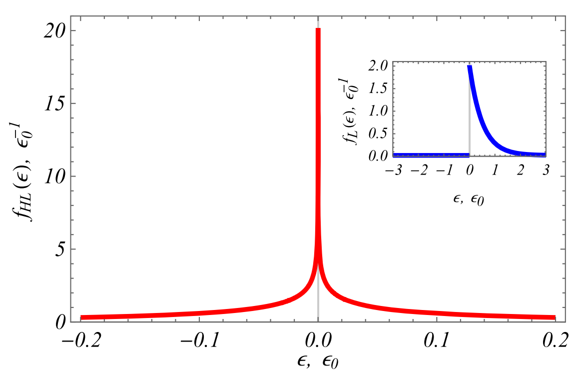


FIG. 1. (Main panel) Energy distribution function $f_{\mathrm{HL}}(\epsilon)$ of a half-leviton excited out of a zero-temperature Fermi sea with the help of a Lorentzian voltage pulse $V(t)$ carrying half of the flux quantum, $(e/\hbar)\int dtV(t) = \pi$. The energy $\epsilon = E - \mu$ is counted from the Fermi energy $\mu$ and is normalized to $\epsilon_0 = \hbar/\Gamma_\tau$, with $\Gamma_\tau$ being the half-width of a voltage pulse. The peak at zero energy is $f_{\mathrm{HL}}(\epsilon \to 0) \approx (2/\pi^2\epsilon_0)\ln^2(\epsilon_0/|\epsilon|)$. (Inset) Energy distribution function of a leviton, a particle with an integer charge $e$ excited by a voltage pulse $2V(t)$: $f_L(\epsilon > 0) = (2/\epsilon_0)\exp(-2\epsilon/\epsilon_0)$ and $f_L(\epsilon < 0) = 0$ [33].

 

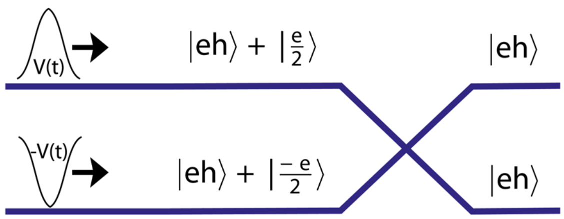


FIG. 2. A sketch of an electronic wave splitter with colliding states excited by the Lorentzian voltage pulses of opposite signs, $V(t)$ and $-V(t)$, carrying half of the flux quantum each. One state is composed of electron-hole pairs, $|eh\rangle$, and a half-leviton, $|e/2\rangle$, and the other one is composed of electron-hole pairs and an anti–half-leviton, $|-e/2\rangle$. Only electron-hole pairs contribute to the excess correlation function of the state projected onto one of the outputs of a symmetric wave splitter.

coherent annihilation on a wave splitter is impossible with ordinary quasiparticles, e.g., electrons and holes, whose energies lie above and below the Fermi energy, respectively. Therefore, they can be annihilated only as a result of inelastic processes, which generally break the phase coherence. The elastic collisions of ordinary single electrons and holes do not lead to annihilation [34–37], except for specific setups—for instance, where an electron emitted by one source is passed by and reabsorbed by another source attempting to emit a hole [38]. Another example is a setup where energies of electrons and holes are aligned, but the success rate of annihilation is small [39]. A possibility for an effective coherent annihilation of particles on a wave splitter predicted here opens a route for entangling Fock states with different numbers of fermions in solid-state quantum circuits.

*Half-leviton.*—To characterize quasiparticles arising in a one-dimensional chiral or ballistic system of noninteracting spinless electrons under the action of a dynamic source, we introduce the excess first-order correlation function [40,41]. This function is defined as the difference of electronic correlation functions with the source on and off, $G^{(1)}(1;2) = \langle\hat{\Psi}^\dagger(1)\hat{\Psi}(2)\rangle_{\text{on}} - \langle\hat{\Psi}^\dagger(1)\hat{\Psi}(2)\rangle_{\text{off}}$. Here, $\Psi(j)$ is an electron field operator calculated at point $x_j$ and time $t_j$ behind the source. The quantum statistical average $\langle\ldots\rangle$ is taken over the equilibrium state of an electron system incoming to the place where the source is located. Incoming electrons are described by the Fermi distribution function with the chemical potential $\mu$ and temperature $\theta$. We will utilize the wideband approximation, where all of the relevant energy scales are small compared to $\mu$ and the spectrum of electrons of the Fermi sea can be linearized around the Fermi energy. In such a case, the excess correlation function depends on a reduced time $t_j \equiv t_j - x_j/v_\mu$ (with $v_\mu$ being the Fermi velocity), rather than on space and time coordinates separately.

If quasiparticles are excited by a time-dependent voltage $V(t)$, then, at zero temperature, the excess correlation function is [42]

$$G^{(1)}(t_1;t_2) = \frac{e^{i(t_1-t_2)(\mu/\hbar)}}{v_\mu}\frac{e^{i(e/\hbar)\int_{t_2}^{t_1}dt' V(t')} - 1}{2\pi i(t_1-t_2)}. \quad (1)$$

Here, we are interested in a Lorentzian voltage pulse of width $2\Gamma_\tau$, $eV(t) = n^\star 2\hbar\Gamma_\tau/(t^2+\Gamma_\tau^2)$, which carries a flux $\varphi = 2\pi n^\star$. The corresponding correlation function is denoted as $G^{(1)}_{n^*}$.

For $n^\star = 0.5$, the exponential function can be expressed in terms of algebraic functions, which allows us to represent $G^{(1)}_{0.5}$ as the sum of two terms,

$$G^{(1)}_{0.5}(t_1;t_2) = \frac{e^{i(t_1-t_2)(\mu/\hbar)}}{v_\mu}\{g_{\text{HL}}(t_1;t_2) + g^{(1)}_{eh}(t_1;t_2)\},$$
$$g_{\text{HL}}(t_1;t_2) = \frac{\Gamma_\tau}{2\pi\sqrt{t_1^2+\Gamma_\tau^2}\sqrt{t_2^2+\Gamma_\tau^2}},$$
$$g^{(1)}_{eh}(t_1;t_2) = \frac{\frac{t_1t_2+\Gamma_\tau^2}{\sqrt{t_1^2+\Gamma_\tau^2}\sqrt{t_2^2+\Gamma_\tau^2}} - 1}{2\pi i(t_1-t_2)}. \quad (2)$$

The first term, $g_{\text{HL}}$, is factorized into the product of two terms dependent on a single time each. It describes a single-particle excitation since all corresponding higher-order correlation functions are identically zero [43]. I call it a half-leviton because it is excited by a half voltage pulse, which excites a leviton [7], and mark corresponding quantities by a subscript HL. This excitation carries a charge $q^\star_{\text{HL}} = e\int dt g_{\text{HL}}(t;t) = e/2$. To understand why a charge is fractional, one needs to note that the state of HL is a mixed state. This follows from the fact that the purity coefficient [42] calculated for $g_{\text{HL}}$ is less then one: $P_{\text{HL}} = \int dt g_{\text{HL}}(t_1;t) g_{\text{HL}}(t;t_2)/g_{\text{HL}}(t_1;t_2) = 0.5$. Since $q^\star_{\text{HL}} = eP_{\text{HL}}$, one can say that the state in question corresponds to a single particle with an integer charge $e$ appearing with probability $P_{\text{HL}}$ and a vacuum state appearing with probability $1-P_{\text{HL}}$. Therefore, $q^\star_{\text{HL}}$ is an effective charge.

Using the purity coefficient, we can write, $g_{\text{HL}}(t_1;t_2) = P_{\text{HL}}\Phi^*_{\text{HL}}(t_1)\Phi_{\text{HL}}(t_2)$, and find that a corresponding single-particle wave function $\Phi_{\text{HL}}(t)$ can be chosen to be real valued,

$$\Phi_{\text{HL}}(t) = \sqrt{\frac{\Gamma_\tau}{\pi}}\frac{1}{\sqrt{t^2+\Gamma_\tau^2}}, \quad (3)$$

and normalized to one, $\int dt|\Phi_{\text{HL}}(t)|^2 = 1$. Note that this wave function is symmetric in time, $\Phi_{\text{HL}}(t) = \Phi_{\text{HL}}(-t)$.

The second term in Eq. (2) describes electron-hole excitations (hence the subscript $eh$), which do not carry any charge, $I_{eh}(t) \equiv eg^{(1)}_{eh}(t;t) = 0$. Their presence can be verified via the shot noise measurement [7,19] or with the help of an interference current [44,45]). The electron-hole pair state is a multiparticle state whose single-particle components are orthogonal to a half-leviton state.

Electron-hole pairs do carry energy injected by a voltage pulse into an electron system. By contrast, the HL does not carry any energy. To show this, let us go over from time domain to energy domain and introduce the energy distribution function for excited particles (see, e.g., Ref. [46]):

$$f(\epsilon) = \frac{v_\mu}{h} \iint dt_1 dt_2 e^{-i(\mu+\epsilon)[(t_1-t_2)/\hbar]} G^{(1)}(t_1;t_2), \quad (4)$$

where $\epsilon$ is an energy counted from the Fermi energy. The function $f(\epsilon)$ is a probability density for finding an excited particle with energy $\epsilon$. Using a correlation function given in Eq. (2), we find

$$f(\epsilon) = f_{\mathrm{HL}}(\epsilon) + f_{eh}(\epsilon),$$

$$f_{\mathrm{HL}}(\epsilon) = \frac{P_{\mathrm{HL}}}{h} \left| \int dt \cos(\epsilon t/\hbar) \Phi_{\mathrm{HL}}(t) \right|^2,$$

$$f_{eh}(\epsilon) = \frac{1}{h} \iint dt_1 dt_2 \sin(\epsilon[t_2 - t_1]/\hbar) i g_{eh}^{(1)}(t_1;t_2). \quad (5)$$

In the last equation, I used $g_{eh}^{(1)}(-t_1;-t_2) = g_{eh}^{(1)}(t_1;t_2)$.

The distribution function is normalized such that $\int d\epsilon f(\epsilon) = P_{\mathrm{HL}}$. Electron-hole pairs do not contribute to this equation. The reason is as follows. By definition, electron and hole contributions to the excess correlation function—and, correspondingly, to the distribution function $f_{eh}(\epsilon)$—have opposite signs. As a result, $\int d\epsilon f_{eh}(\epsilon) = 0$. To get the number of either electrons or holes separately, we have to integrate $f_{eh}(\epsilon)$ over either positive or negative energies only.

The distribution function of a half-leviton is shown in Fig. 1. It is an even function of energy, $f_{\mathrm{HL}}(\epsilon) = f_{\mathrm{HL}}(-\epsilon)$, and therefore it does not contribute to the energy of excitations $\langle\epsilon\rangle = \langle\epsilon\rangle_{\mathrm{HL}} + \langle\epsilon\rangle_{eh}$,

$$\langle\epsilon\rangle_{\mathrm{HL}} = \int d\epsilon \epsilon f_{\mathrm{HL}}(\epsilon) = 0. \quad (6)$$

This is why I call a half-leviton a *zero-energy excitation*. In contrast, a true leviton, excited by a voltage pulse with $n^\star = 1$, has a nonzero energy, $\langle\epsilon\rangle_L = \int d\epsilon f_L(\epsilon)\epsilon = \hbar/(2\Gamma_\tau)$ [33]; see the inset in Fig. 1 for the leviton's distribution function, $f_L(\epsilon)$.

Note that the HL's energy is zero on average, but it does fluctuate. This fact differentiates HL from quasiparticles in Majorana zero modes in topological insulators and superconductors, whose energy is strictly zero; see, e.g., Refs. [47,48]. In addition, HL is charged while a Majorana fermion is neutral.

The distribution function for electron-hole pairs is an odd function of energy, $f_{eh}(\epsilon) = -f_{eh}(-\epsilon)$. Therefore, electron-hole pairs do carry (excess) energy, which is pumped by a time-dependent voltage $V(t)$ into the Fermi sea:

$$\langle\epsilon\rangle_{eh} \equiv \int d\epsilon f_{eh}(\epsilon)\epsilon = i\hbar \int dt \left.\frac{\partial g_{eh}^{(1)}(t;t')}{\partial t'}\right|_{t'=t} = \frac{1}{4}\frac{\hbar}{2\Gamma_\tau}. \quad (7)$$

This energy is a quarter of the energy of a leviton.

The same result follows also from a time-dependent heat current, $J_Q(t)$, induced by a voltage pulse [49]. At zero temperature, one can find quite generally [50] that a charge current $I(t) = eg(t;t)$ and a heat current, both induced by a voltage pulse in a single-channel chiral or ballistic conductor, comply with the Joule law,

$$J_Q(t) = R_q I^2(t), \quad (8)$$

where $R_q = h/(2e^2)$ is the charge relaxation resistance [51], the Büttiker resistance [52]. Heat is nothing but the excess energy carried by the excitations [53]. Indeed, $\int dt J_Q(t) = \hbar/(8\Gamma_\tau)$, which agrees with Eq. (7).

The fact that the Joule law, Eq. (8), works in the present case is remarkable since charge and heat are carried by different pieces of the excited state, the HL and electron-hole pairs, respectively. Actually, these pieces can be separated in experiment. To show this, let us first consider what is excited by a voltage pulse of the opposite sign.

*Anti–half-leviton.*—In the case of $n^\star = -0.5$, the correlation function is $G_{-0.5}^{(1)}(t_1;t_2) = e^{i(t_1-t_2)(\mu/\hbar)}\{-g_{\mathrm{HL}}(t_1;t_2) + g_{eh}^{(1)}(t_1;t_2)\}/v_\mu$. The change of voltage sign does not alter the electron-hole part. What is changed is the sign of a charge of a single-electron excitation, which I will call an *anti–half-leviton* (aHL).

Let us take states with HLs and aHLs that are excited at different contacts and mix them at a wave splitter, a quantum point contact, with transmission $T$ and reflection $R = 1 - T$ being the probabilities. The correlation function of excitations at output is $G_{\mathrm{out}}^{(1)} = TG_{0.5}^{(1)} + RG_{-0.5}^{(1)}$. In the case of a symmetric wave splitter, $T = R = 0.5$, we find $G_{\mathrm{out}}^{(1)}(t_1;t_2) = e^{i(t_1-t_2)(\mu/\hbar)} g_{eh}^{(1)}(t_1;t_2)/v_\mu$. That is, the correlation function of the state projected onto the output channel shows the presence only of electron-hole pairs; see Fig. 2. On the level of the electronic correlation function, the HL and the anti-HL effectively annihilate each other and the projected outgoing state contains only electron-hole pairs. The measurement of a correlation function made on such a state can serve as the reference point for a measurement made on $G_{0.5}^{(1)}$ in order to extract characteristics of a half-leviton.

Note that the excess first-order correlation function contains all information about excitations: their charge, energy, fluctuations, coherence times, etc. The correlation function is additive and, therefore, it is specifically suitable for the electron-hole pair elimination procedure outlined above. The correlation function can be directly measured with the help of an interference current as was suggested in

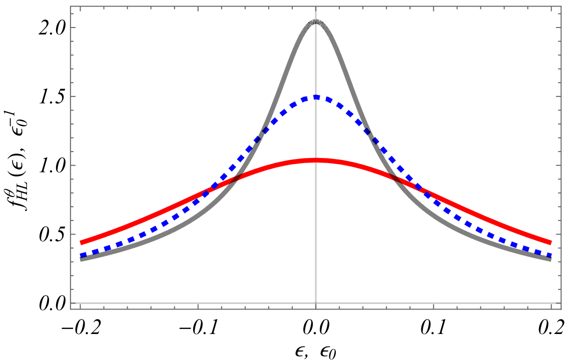


FIG. 3. Energy distribution function $f^{\theta}_{\mathrm{HL}}(\epsilon)$, Eq. (9), of a half-leviton at different temperatures of the Fermi sea $\theta = \theta^{\star}/k$, where $\theta^{\star} = \hbar/(\pi k_B \Gamma_\tau)$ and $k = 5$ (the red solid line), $k = 10$ (the blue dashed line), and $k = 20$ (the black short-dashed line). For a voltage pulse with the width $2\Gamma_\tau = 30$ ps [7], the characteristic temperature is $\theta^{\star} \approx 160$ mK.

Refs. [41,54]. Moreover, the distribution function $f_{\mathrm{HL}}(\epsilon)$ (Fig. 1) can be measured using already available experimental tools, quantum dots as energy filters [55], or the shot noise spectroscopy [7,23,56]. One should note that a finite temperature of the Fermi sea restricts the precision of measurement of a zero-energy peak.

*Temperature effect.*—At a nonzero temperature, $\theta > 0$, the excess correlation function is $G^{(1)}_{0.5,\theta}(t_1; t_2) = \eta[(t_1 - t_2)/\tau_\theta] G^{(1)}_{0.5}(t_1; t_2)$, where $\eta(x) = x/\sinh(x)$ is a temperature-induced suppression factor and the thermal time $\tau_\theta = \hbar/(\pi k_B \theta)$, with $k_B$ being the Boltzmann constant [42]. Substituting the equation above into Eq. (4) and isolating a part related to the HL, we find

$$f^{\theta}_{\mathrm{HL}}(\epsilon) = \int d\omega \eta_\omega f_{\mathrm{HL}}(\epsilon + \hbar\omega), \tag{9}$$

where $\eta_\omega = (\pi\tau_\theta/4)\cosh^{-2}(\pi\omega\tau_\theta/2)$ is the Fourier transform of $\eta(t/\tau_\theta) = \int d\omega e^{-i\omega t}\eta_\omega$. A zero-temperature distribution function $f_{\mathrm{HL}}(\epsilon)$ is given in Eq. (5). The function $f^{\theta}_{\mathrm{HL}}(\epsilon)$ is presented in Fig. 3 for different temperatures for the Fermi sea. Though a zero-energy peak is suppressed with increasing temperature, its shape remains symmetric around $\epsilon = 0$.

*Conclusion.*—A dynamically perturbed Fermi sea can host exotic zero-energy excitations with an effectively fractional charge. In this Letter, I discuss an example of such a quasiparticle, which can be excited by a Lorentzian voltage pulse $V(t)$ with a half-integer flux, $\varphi = (e/\hbar)\int dt V(t) = \pi$, using the same technique that is used to generate levitons [7,23]. A single particle with an effective charge $e/2$, a HL, is excited together with a cloud of electron-hole pairs, which, however, can be isolated and used as a reference point for studying the HL. A half-leviton is described by a single-particle state, which is mixed in equal proportions with the vacuum state—hence, a fractional charge. This single-particle state is a coherent superposition of states with energies symmetrically placed near the Fermi energy. Therefore, the energy of the HL counted from the Fermi energy is zero. The wave function of the HL is real valued and, therefore, it remains the same when we go over to an anti-HL, a particle excited by a voltage pulse with $\varphi = -\pi$. These properties enable the HL and the anti-HL to effectively annihilate each other while colliding at an electronic wave splitter, what paves the way for entangling fermionic Fock states with a different number of particles. Dynamic excitation of an electron many-particle system is an exciting and promising platform for quantum coherent electronics, which "… does not require delicate nanolithography, considerably simplifying the circuitry for scalability" [7]. Moreover, hunting for new types of excitations simply means reshaping a voltage pulse.

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