Positive Magnetoresistance of Composite Fermions in Laterally Modulated Structures

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(March 24, 2022)

Adopting the mean-field composite fermion picture, we describe the magneto-transport properties of a two-dimensional electron gas with laterally modulated density around filling factor $\nu = 1/2$. The occurrence of a strong positive magnetoresistance at low effective magnetic fields as well as Weiss oscillations, which were observed in recent experiments in such systems, can be explained within a semi-classical Boltzmann equation approach, provided one goes beyond a second order approximation in the modulation strength.

PACS numbers: 71.10.Pm, 73.40.Hm, 73.20.Dx

Transport properties of strongly correlated two-dimensional electron systems (2DES) in high magnetic fields, where the lowest Landau level is about half filled, seem to be surprisingly well described by the “Composite Fermion” (CF) picture. This model is motivated by a singular Chern-Simons gauge-field transformation that maps the electron system near filling factor $\nu = 1/2$ onto a metal of quasi-particles (the CFs). They can be interpreted as electrons to which two magnetic flux quanta have been attached [1]. In the usual mean-field approximation, the CFs are considered as having the same density $n_{\text{CF}} = n_{\text{el}} \equiv n$ as the electrons and as moving in the weak effective magnetic field $B_{\text{eff}} = B - 2\Phi_0 n$, where $B$ is the external applied magnetic flux density and $\Phi_0 = h/e$ the flux quantum. Halperin, Lee, and Read [4] predicted that, at low temperatures, these CFs fill in momentum space the spin-polarized states within a well defined Fermi circle of radius $k_F = \sqrt{4\pi n}$. Therefore, one expects that transport properties, which are governed by elastic scattering at the Fermi energy, can be calculated without considering explicitly the mutual interaction between the CFs, by analogy with the Landau theory of the electron Fermi liquid in zero and weak magnetic fields. Indeed, interesting features of the magnetoresistance of 2DESs (near $\nu = 1/2$) in lateral superlattices, such as two-dimensional anti-dot lattices or one-dimensional (1D) superlattices, created either dynamically by the application of surface acoustic waves or statically by surface etching, have been reproduced astonishingly well by calculations describing the CFs as classical, non-interacting particles moving in suitable effective fields [3, 8].

Here, we focus on static 1D density modulations and compare experimental de-transport measurements [8] with calculations based on the linearized Boltzmann equation (LBE). The purpose of the present letter is to demonstrate that all characteristic features of the magnetoresistance curve near $B = B_{1/2} \equiv 2\Phi_0 n$ can be reproduced in the quasi-classical CF picture, provided one incorporates all modulating fields and solves the LBE beyond the Beenakker-type approximation (BA) [6], which allows an analytical solution and has been employed in previous work [3, 8]. These features are: (i) a pronounced V-shaped minimum near $B_{1/2}$, explained by “channeled orbits” which are omitted in the BA, (ii) shoulders or minima related to commensurability effects (Weiss oscillations due to drifting cyclotron orbits), (iii) a steep increase of the resistance with $|B_{\text{eff}}|$ between the structures caused by commensurability effects and the oscillatory structures due to the fractional quantum Hall effect, and (iv) an asymmetry in the slope and magnitude of these steep resistance flanks with respect to $B_{\text{eff}} = 0$, due to interference effects between the direct electrostatic density modulation and the induced modulation of $B_{\text{eff}}$. The material parameters needed to achieve agreement with the experiment around $\nu = 1/2$ are consistent with those needed to explain the commensurability effects at low magnetic fields ($B < 0.5$ T).

Recent work by Mirlin et al. [8], based on an analytical solution of the LBE within the BA for a special model of anisotropic scattering [8], obtained reasonable results for the features (ii) and (iii), but not for (i) and (iv). Their approximation misses relevant physics at low $B_{\text{eff}}$, where the CF picture is expected to apply best, and gives a rather poor fit to the experiment [8] for very small $|B_{\text{eff}}|$ (see Fig. 1 below).

To describe the resistance for $B < 0.5$ T, we follow Ref. 4 for pure electric modulation, and treat the 2DES as a degenerate Fermi gas of non-interacting particles, with charge $-e$ and effective mass $m_{\text{el}}^*$, average density $\bar{n}$, Fermi energy $E_{\text{F}}^{\text{el}} = \bar{n}/D_0^{\text{el}}$, and density of states $D_0^{\text{el}} = m_{\text{el}}^*/(\pi \hbar^2)$. The etching of grooves into the surface of GaAs-heterostructures is assumed to produce an external electrostatic potential energy $V^{\text{ext}}(x) = V_0^{\text{ext}} \cos qx$ of period $a = 2\pi/q$ in the plane of the 2DES. It is screened by the 2DES and this leads within the Thomas-Fermi approximation to a potential energy $V^{\text{el}}(x) = V_0^{\text{rel}} \cos qx$ with $V_0^{\text{rel}} = V_0^{\text{ext}}/[1 + 2/(qa_{\text{Bohr}})]$, where $a_{\text{Bohr}} = 1/(\pi e^2 D_0^{\text{el}}/\hbar^2)$ is the effective Bohr radius. For GaAs, $\kappa = 12.4$ and $a_{\text{Bohr}} \approx 10$ nm $< a \sim 400$ nm. Consequently, the relative modulation strength is $\epsilon_{\text{el}} = V_0^{\text{rel}}/E_{\text{F}}^{\text{el}} \approx V_0^{\text{ext}} \kappa a/(e^2 \bar{n})$, and the modulated electron density is $n(x) = \bar{n}[1 - \epsilon_{\text{el}} \cos qx]$, independent of $m_{\text{el}}^*$.

To describe the resistance of CFs, we modify the ap-
proach of Ref. [9]. We treat the CF system as a degenerate Fermi gas of non-interacting particles, with charge $-e$, effective mass $m_{\text{CF}}$, average density $\bar{n}$, Fermi energy $E_{\text{CF}}^F = \bar{n}/D_{0}^F$, and density of states $D_{0}^F = m_{\text{CF}}/(2\pi\hbar^2)$, which obeys Newton’s equation, $\bar{m}_{\text{CF}} \ddot{v} = -e[F_{\text{eff}} + v \times (B_{\text{eff}} \omega_z)]$, with effective electric and magnetic fields. We use the LBE to calculate the response to an external homogeneous electric field $E^{(0)}$. In the absence of $E^{(0)}$, $F_{\text{eff}} = \nabla V(x)/e$ is determined by the screened modulation potential, which we parametrize as $V(x) = \epsilon_{\text{CF}} E_{\text{CF}}^F \cos(qx)$. Since the static Thomas-Fermi screening of the CFs should be equivalent to that of electrons (with $D_{0}^F > D_{0}^e$ instead of $D_{0}^e$), we expect $\epsilon_{\text{CF}} \approx \epsilon_{e}$, so that the modulated density $n(x) = \bar{n}[1 - \epsilon_{\text{CF}} \cos qx]$ for $B$ near $B_{1/2}$ is the same as that for small $B$ within this approximation. This yields the effective magnetic field $B_{\text{eff}}(x) = B - 2\Phi_{0} m(x) = B_{\text{eff}} + \delta B_{\text{eff}}(x)$, where $\delta B_{\text{eff}}(x) = \epsilon_{\text{CF}} B_{1/2} \cos qx$ is the magnetic modulation.

The linear response to $E^{(0)}$ is carried by CFs at the Fermi edge, $(m_{\text{CF}}^* / 2v^2) + V(x) = E_{\text{CF}}^F$. Due to the translational invariance in $y$ direction, the (suitably scaled) zero-temperature distribution function $\Phi(x, \varphi)$ depends only on two variables, the position $x$ and the polar angle $\varphi$ of the velocity $v(x, \varphi) = v(x) (\cos \varphi, \sin \varphi)$, with $v(x) = v_F [1 - V(x)/E_{\text{CF}}^F]^{1/2}$, $v_F = \hbar k_{\text{CF}}^F / m_{\text{CF}}$, and $k_{\text{CF}}^F = [4\pi \bar{n}]^{1/2}$. We normalize $\Phi(x, \varphi)$ so that the current density $j(x)$ reduces to

$$j(x) = e^2 D_0^F \int_{-\pi}^{\pi} \frac{d\varphi}{2\pi} v(x, \varphi) \Phi(x, \varphi).$$

A current of CFs implies motion of flux tubes and produces the Chern-Simons electric field $E_{\text{CS}}$

$$E_{\text{CS}} = -(2\hbar/e^2) j \times \mathbf{e}_z,$$

which, in the linearized Boltzmann equation,

$$\mathcal{D} \Phi - C[\Phi] = v(x, \varphi) \left( E^{(0)} + E_{\text{CS}}(x) \right),$$

adds to the external driving field $E^{(0)}$. Therefore a self-consistent solution of Eqs. (1)-(3) is necessary. The drift operator

$$\mathcal{D} = v(x) \cos \varphi \partial_x + (\omega_c + \omega_w(x) + \omega_e(x, \varphi)) \partial_\varphi,$$

contains the average value and the periodic component of the effective magnetic field via $\omega_c = eB_{\text{eff}}/m_{\text{CF}}$, and $\omega_w(x) = eB_{\text{eff}}(x)/m_{\text{CF}}$, respectively. The electric modulation enters in $\omega_e(x, \varphi) = -\sin \varphi \partial v/dx$ and $v(x)$ [3]. The collision operator is taken to be the same as for the unmodulated system and is written in the form

$$C[\Phi] = \frac{1}{\tau} \int_{-\pi}^{\pi} \frac{d\varphi'}{2\pi} P(\varphi' - \varphi) \left[ \Phi(x, \varphi') - \Phi(x, \varphi) \right],$$

where the differential cross section is parametrized by a relaxation time $\tau$ and a dimensionless kernel $P(\varphi)$. For actual calculations we take $P(\varphi) = b + (1 - b) P_p(\varphi)$ with $0 \leq b \leq 1$ and

$$P_p(\varphi) = [(2p^2 p)/((2p)!)] \cos^{2p}(\varphi/2),$$

which has the finite Fourier expansion $P(\varphi) = \sum_{n=0}^{p} \gamma_n \cos(n \varphi)$ with $\gamma_0 = 1$ and $\gamma_n = 2(1 - b)(p!)^{2} / ((p + n)! (p - n)!)$ for $n \geq 1$. For $b = 1$ and $p = 0$, $P(\varphi)$ describes isotropic scattering. For $b < 1$ the fraction $(1 - b)$ of the total scattering cross section is due to anisotropic scattering. With increasing $p$, $P_p(\varphi)$ is increasingly stronger peaked in forward direction.

To solve Eqs. (1)-(3), we calculate the distribution function of the homogeneous, unmodulated CF system first. This yields the resistivity tensor $\rho_{\mu \nu}$ with components given by $\rho_{xx}^\mu = \rho_{yy}^\mu = \rho_{0}^\mu$ and $\rho_{xy}^\mu = -\rho_{yx} = \omega_{1/2} \tau_{1} \rho_{0}^\mu$, where $\rho_{0}^\mu = m_{\text{CF}}^\mu / (e^2 \hbar \lambda_{\text{CF}})$, and $\tau_{1} = \tau / (1 - \gamma_1/2)$ is the relevant CF transport scattering time. Although the CFs move in the effective magnetic field $B_1/2$, the Hall resistance $\rho_{xy}^\mu$ is determined by the cyclotron frequency $\omega_{1/2} = \omega + \omega_{1/2}$ for the total magnetic field $B$, with $\omega_{1/2} = eB_1/\mu_{\text{CF}}$ due to the inclusion of the Chern-Simons electric field. Introducing the mean free path $\lambda_{\text{CF}} = \tau \hbar k_{\text{CF}}^F / m_{\text{CF}}^\mu$, we obtain $\rho_{0}^\mu = \hbar k_{\text{CF}}^F / (e^2 \hbar \lambda_{\text{CF}})$ and $\rho_{xy}^\mu = B/(ne)$, so that the resistivity tensor is independent of $m_{\text{CF}}^\mu$, and the Hall resistance is the same as that of the 2DES. Comparing with the corresponding Drude formulas for the 2DES at low $B$, we see that the ratio of the longitudinal resistances, $\rho_{xx}^\mu / \rho_{yy}^\mu = \sqrt{2} \lambda_{\text{CF}} / \lambda_{\text{CF}}$, directly reflects the ratio of the corresponding mean free paths.

The macroscopic resistivity tensor $\rho$ of the modulated system, relates the spatial average $\langle j(x) \rangle$ of the current density $j(x)$ to the driving field, $\rho \langle j(x) \rangle = E^{(0)}$. Exploiting the solution for the homogeneous system and the continuity equation, which implies that $j_{x}(x)$ for the modulated system is independent of $x$, it can be shown that $\rho$ differs from $\rho^0$ only in its $x$-component. With some formal but exact manipulations (similar to those of ref. [9]) one finds

$$\langle g_{xx}^\mu - \rho_{yy}^\mu \rangle / \rho_{0}^\mu = G / [1 - G / (1 + (\omega_{1/2} \tau_{1}^2))^2].$$

In this equation

$$G = \frac{2\tau_1}{v_F} \int_{-\pi}^{\pi} \frac{d\varphi}{2\pi} g(x, \varphi) \chi(x, \varphi),$$

where $g(x, \varphi) = v_F^2 (dV/dx) / (2E_{\text{CF}}^F) + \omega_{1/2} v_{y}$ is of first order in the modulation strength $\epsilon_{\text{CF}}$, and $\chi(x, \varphi)$ is the solution of the modified Boltzmann equation

$$\mathcal{D} \chi - C[\chi] = g(x, \varphi) + \omega_{1/2} v_y (j_{x}^1 - j_{y} - 1).$$

The basic differences between these results for the CF system and those for an electron system [3] derive from inclusion of the Chern-Simons electric field in Eq. (3):
It is responsible for the occurrence of $\omega_{\text{tot}} \propto B$ in the denominator of Eq. (7), instead of $\omega \propto B - B_{1/2}$, and the last term in the right hand side of Eq. (8). This last term can be expressed in terms of $G$, $\tau_{\text{tr}}$, and $\omega_{\text{tot}} \tau_{\text{tr}}^2$, and requires a self-consistent solution of Eqs. (7)-(9). Inserting $\tau_{\text{tr}} = \lambda_{\text{CF}}/v_{\text{CF}}$ into Eq. (8) makes $G$ and $\bar{g}_{xx}$ independent of $m_{\text{CF}}$, as it should be.

Numerical solution of Eqs. (7)-(9) for the CF system and of the corresponding equations of Ref. [9] for the electron system allows to calculate the resistance of the modulated sample for strong applied fields $B \approx B_{1/2}$ and in the low $B$ regime, respectively. To compare with experiment, we need to choose a reasonable set of parameters. Whereas in both $B$ regimes the average density $\bar{n}$ and the modulation period $a$ are the same, and also the modulation strengths $\epsilon_{\text{CF}} \approx \epsilon_{\text{el}}$ should be similar, the parameters characterizing the collision operator (in our model $\tau$, $b$ and $p$) are expected to be different. To reproduce the characteristics of the Weiss oscillations observed at small $B \ll B_{1/2}$, we have to assume a very anisotropic differential cross section with a sharp peak in forward direction and a zero-$B$ resistance of only a few $\Omega$ ($\lambda_{\text{el}} \sim 50 \mu\text{m}$). This is reasonable since these high-mobility samples have a spacer thickness of about 60 nm, so that the donors produce a smooth random potential in the plane of the 2DES that scatters the electrons predominantly under small angles. The CFs, on the other hand, are scattered not only by these small-amplitude donor-induced random potential fluctuation, but also – and predominantly – by the large-amplitude random fluctuations of the effective magnetic field that originates from the donor-induced density fluctuation. This results in $\lambda_{\text{CF}} \ll \lambda_{\text{el}}$ and a less pronounced forward scattering for the CFs.

Figure 1 shows resistivity data for a sample with density $\bar{n} = 1.82 \cdot 10^{11} \text{ cm}^{-2}$ and period $a = 400 \text{ nm}$ near filling factor $\nu = 1/2$ [7]. The thick short-dashed line is the experiment of Ref. [8]. The thin dash-dotted line is the theoretical curve using the approximations of Ref. [8], based on the Beenakker-type approximation (BA) mentioned above. It neglects the modulation effects in the drift operator $D$, that causes a resistance correction quadratic in $\epsilon_{\text{CF}}$. Furthermore, the direct effect of the electric modulation on the CFs was ignored in this fits. Neglecting this direct effect and using the BA, we can closely reproduce this symmetric curve of Ref. [8] using the same values for $\bar{n}$, $\epsilon_{\text{CF}}$, and using $b = 0$, $p = 6$ for our model of the scattering cross section, Eqs. (7) and (8). If we include the direct electric modulation in addition to the magnetic one, we obtain the asymmetric thin dotted line. For small $|B_{\text{eff}}|$, this again closely reproduces the results of Ref. [8], but at large $|B_{\text{eff}}|$ the curve shows a distinct asymmetry, similar to the experimental curve, and similar to the asymmetry obtained for mixed electric and magnetic modulations in 2DESs at low $B$ [9,10]. Solved Eqs. (7)-(9) for the same set of parameters but without any approximation of the drift opera-
tor, we obtain the thin solid curve of Fig. 1. This shows that approximating the drift operator $D$ by the one of the homogeneous system is insufficient for $|B_{\text{eff}}| < 0.5 \text{ T}$. Physically the BA means that the forces due to the modulation fields are neglected against the Lorentz force due to the average magnetic field $B_{\text{eff}}$. This omits the effect of channeled orbits, i.e. of “snake orbits” [3] along the lines of vanishing $B_{\text{eff}}(x)$ which, similar to the “open orbits” [11] in the case of a pure electric modulation, cause a pronounced positive low-field magneto-resistance [12]. Since $B_{1/2} = 15 \text{ T}$ [10], the effective magnetic modulation has an amplitude $\epsilon_{\text{CF}} |B_{1/2} - 0.45 \text{ T} |$, so that the effect of channeled trajectories should be important for $|B_{\text{eff}}| < 0.45 \text{ T}$, and a second order approximation in $\epsilon_{\text{CF}}$ is not adequate for the $B_{\text{eff}}$ values shown in Fig. 1.

Comparing the thin solid curve with the experimental one suggests that the pronounced V-shaped resistance minimum at $B_{\text{eff}} = 0$ seen in the experiment is indeed due to CFs on channeled orbits. To achieve a better quantitative agreement, we choose more realistic parameters. First we take $b = 0.1$ and $p = 2$, since then our model [8] approximates closely (apart from the $\varphi = 0$ divergence) the differential cross section derived by Aronov et al. [15] for random magnetic field scattering. This choice allows a good fit of the resistance near $\nu = 1/2$ for all samples available to us, indicating the same scattering mechanism in those samples. Then, assuming a density modulation of about 3%, we have to take $\epsilon_{\text{CF}}^0 = 650 \text{ } \Omega$ ($\lambda_{\text{CF}} \approx 1.3 \text{ a}$), i.e. a larger value than the resistance of the unmodulated reference sample ($\epsilon_{\text{CF}}^0 = 270 \text{ } \Omega$) taken in Ref. [8]. This yields the thick solid curve from the exact solution of the LBE, which reproduces nicely all features of the experimental data. The larger $\epsilon_{\text{CF}}^0$ value is reasonable since the etching procedure, in addition to the intended modulation, introduces unintended defects which increase the resistance.

For the sample discussed so far ($a = 400 \text{ nm}$), the mean free path of the CFs is only marginally larger than the modulation period, $\lambda_{\text{CF}} \approx 1.3a$, so that commensurability effects are visible only as weak shoulders (thick lines of Fig. 1). To see these effects more clearly, we have investigated a sample with a smaller modulation period, $a = 275 \text{ nm}$, so that we expect $\lambda_{\text{CF}} \approx 1.9a$. Experimental [14] and theoretical results (with the same anisotropy parameters $b$ and $p$ as in Fig. 1) are plotted in Fig. 2. Indeed, the commensurability effects near $|B_{\text{eff}}| = 0.5 \text{ T}$ are more pronounced. Again, the low field resistance (inset of Fig. 2) and that near $\nu = 1/2$ can be reproduced theoretically with a consistent set of parameters.

In conclusion, our numerical solution of the linearized Boltzmann equation for composite Fermions without the Beenakker-type approximation used in previous work [3,8], and thus the inclusion of the effect of “channeled orbits”, reproduces the typical V-shaped resistance minimum at $\nu = 1/2$ ($\tilde{B}_{\text{eff}} = 0$) and yields good overall agreement with experiment, consistent with the low-field
magnetoresistance oscillations. Including further the superposition of the original electric superlattice and the induced effective magnetic superlattice, we can also reproduce the asymmetry of the resistance curve around $B_{\text{eff}} = 0$.

We are grateful to J. Smet for providing the unpublished data of Fig. 2 and for helpful discussions. This work was supported by BMBF Grant No. 01BM622.

[1] J. K. Jain, Phys. Rev. Lett. 63, 199 (1989).
[2] B. I. Halperin, P. A. Lee, and N. Read, Phys. Rev. B 47, 7312 (1993).
[3] R. Fleischmann, T. Geisel, C. Holzknecht, and R. Ketzmerick, Europhys. Lett. 36, 167 (1996).
[4] A. D. Mirlin and P. Wölfle, Phys. Rev. Lett. 78, 3717 (1997).
[5] F. von Oppen, A. Stern, and B. I. Halperin, Phys. Rev. Lett. 80, 4494 (1998).
[6] J. Smet, K. von Klitzing, D. Weiss, and W. Wegscheider, Phys. Rev. Lett. 80, 4538 (1998).
[7] C. W. J. Beenakker, Phys. Rev. Lett. 62, 2020 (1989).
[8] A. D. Mirlin, P. Wölfle, Y. Levinson, and O. Entin-Wohlman, Phys. Rev. Lett. 81, 1070 (1998).
[9] R. Menne and R. R. Gerhardts, Phys. Rev. B 57, 1707 (1998).
[10] S. H. Simon, in Composite Fermions, edited by O. Heinonen (World Scientific, Singapore, 1999), p. 91.
[11] P. D. Ye et al., Phys. Rev. Lett. 74, 3013 (1995).
[12] R. R. Gerhardts, Phys. Rev. B 53, 11064 (1996).
[13] J. E. Müller, Phys. Rev. Lett. 68, 385 (1992).
[14] P. H. Beton et al., Phys. Rev. B 42, 9229 (1990).
[15] A. G. Aronov, A. D. Mirlin, and P. Wölfle, Phys. Rev. B 49, 16609 (1994).
[16] J. H. Smet, private communications
[17] The insert of Fig. 1 shows the measured low-field resistance. We need extremely pronounced forward scattering in order to reproduce the small measured resistance and simultaneously the strong suppression of Weiss oscillations at $B < 0.05$ T. The good agreement is of course limited by the Shubnikov-de Haas oscillations which in these high-mobility samples (at $T = 50$ mK) set in at $B > 0.2$ T.

FIG. 1. Resistivity near filling factor $1/2$, $a = 400$ nm, $\bar{n} = 1.82 \cdot 10^{11}$ cm$^{-2}$; thick dashed line: experiment of Ref. 8, thin dash-dotted line: theory of Ref. 8 (parameters: $\varphi^0 = 270 \Omega$ and $\epsilon_{\text{CF}} = 2.6\%$), thin solid line: results of full calculation (including channeled orbits) with $p = 6$ and $b = 0$, thin dotted line: Beenakker approximation for same parameters, thick solid line: full calculation for $\varphi^0 = 650 \Omega$, $\epsilon_{\text{CF}} = 3.5\%$, $p = 2$ and $b = 0.1$. The inset shows the low-field magnetoresistance; dashed line: experiment, solid line: full calculation based on Ref. 8 for $\varphi^0 = 4.5 \Omega$ and $\epsilon_{\text{CF}} = 2.8\%$, $b = 0$ and $p = 15$.

FIG. 2. Resistivity near filling factor $1/2$, $a = 275$ nm, $\bar{n} = 1.98 \cdot 10^{11}$ cm$^{-2}$; thick dashed line: experimental data, thick solid line: calculation for $\varphi^0 = 700 \Omega$, and $\epsilon_{\text{CF}} = 3.8\%$, cross section parameters $p = 2$ and $b = 0.1$, as in Fig. 1 (thick lines), thin dotted line: Beenakker-type approximation (same parameters). The inset shows the low-field magnetoresistivity; dashed line: experiment, solid line: calculation for $\varphi^0 = 7.5 \Omega$, $\epsilon_{\text{CF}} = 3.6\%$, $b = 0$, and $p = 15$. 
Fig. 1 Zwerschke

\[ \rho_{xx} / \Omega \]

\[ B_{\text{eff}} / T \]

\[ \rho_{xx} / k\Omega \]

\[ 0 \quad 0.1 \quad 0.2 \quad 0.3 \quad B/T \]

\[ -1.5 \quad -1 \quad -0.5 \quad 0 \quad 0.5 \quad 1 \]

\[ \bar{B}_{\text{eff}} / T \]
Figure 2: Zwerschke