The following full text is a publisher's version.

For additional information about this publication click this link.
http://hdl.handle.net/2066/103648

Please be advised that this information was generated on 2018-07-30 and may be subject to change.
I. INTRODUCTION

The recently discovered topological insulators are a matter of intense theoretical and experimental study in contemporary condensed matter physics (for a review of the field, see Ref. 1). The exotic surface states that occur in these materials were confirmed and thoroughly studied using surface sensitive techniques, such as angle-resolved photoemission spectroscopy (ARPES) and scanning tunneling microscopy (STM). These measurements demonstrated the absence of backscattering at the surface, which is believed to be protected by time-reversal symmetry.

Of the three-dimensional (3D) topological insulators, Bi$_2$Se$_3$ has the simplest band structure—a single surface Dirac cone; however, transport experiments on this material present a challenge. The surface conducting state in Bi$_2$Se$_3$ has the simplest band structure—a single surface Dirac cone; however, transport experiments on this material present a challenge. The surface conducting state in Bi$_2$Se$_3$ is difficult to separate from the spurious bulk conductivity, believed to originate from Se vacancies. Even in the lowest carrier-concentration samples metallic bulk behavior is observed. A few successful attempts to surmount this problem were reported, including thinning down the 3D material to decrease the bulk contribution, using gate voltage, pulsed high magnetic fields, and changing the material composition to reduce the bulk conductivity. However, providing conclusive transport evidence of the $\pi$ Berry phase expected from the topology of the surface remains a challenge.

Recent attempts to probe the surface of Bi$_2$Se$_3$ crystals using the Shubnikov–de Haas (SdH) effect resulted in an overwhelming bulk signal. These groups focused on making the samples cleaner in order to increase the bulk resistance. However, the decrease in the bulk carrier concentration led to an increased mobility. Consequently, the 3D oscillations dominated the signal.

In our work we adopt the opposite approach; the SdH effect has an exponential sensitivity on the mobility, whereas the conductance varies only linearly with it. This allows a mobility window in which the SdH effect probes only the protected surface states, albeit the dominance of the bulk over the classical conductance. This can be done by increasing the carrier concentration (thus lowering the bulk mobility), in contrast to the conventional approach.

II. METHODS

Here we report on transport and ARPES measurements in Bi$_2$Se$_3$ flakes. We observe SdH oscillations in the magnetoresistance, and provide evidence that they originate solely from the surface Dirac cone. From these measurements we find the effective mass and carrier concentration and provide a lower bound for the Dingle scattering time. The ARPES data provides evidence of a single Dirac cone and the Fermi wave vector of the surface states obtained is in agreement with the magnetotransport result. We study the evolution of the SdH phase with magnetic fields. The broad field range studied and the pronounced oscillations observed result in unprecedented accuracy at which the frequency can be determined. However, we find that the phase is a strong function of the frequency. The data is consistent with a $\pi$ Berry phase changing with magnetic field. However, the conventional Zeeman effect alone is insufficient to describe our data. Other possibilities, such as zero phase, are discussed.
measurements flakes were cut into narrow bars and measured in a four-contact configuration. Other pieces from the same batch were measured in five-contact and Van der Pauw configurations for a more accurate determination of the resistivity and Hall coefficient. Two Hall probes were used for accurate determination of the perpendicular and parallel field components. ARPES measurements were performed using the He I λ line (hν = 21.2 eV) from a Scienta UV lamp and with an R4000 Scienta analyzer. The crystals for the ARPES measurement were cleaved in situ at vacuum better than 5 × 10⁻¹¹ torr at 20 K.

We took extra care to properly align our voltage leads. Misaligned contacts resulted in a spurious Hall contribution in the longitudinal voltage channel. This contribution manifested itself in an antisymmetric background and a phase shift between positive and negative magnetic fields. Although all samples gave similar frequencies, samples with such Hall contribution were not used for the phase analysis.

### III. RESULTS AND DISCUSSION

In the insets of Fig. 1 we show typical transport properties of a sample with high carrier concentration. The resistivity as a function of temperature for a typical flake is shown in the left inset. A metallic temperature dependence (with a saturation value of ∼280 μΩ cm below 20 K) is observed. The carrier concentration of this flake inferred from the Hall measurements (data are presented in Appendix D) was 8.6 × 10¹⁹ cm⁻³; this gives a bulk mobility of 400 cm² V⁻¹ s⁻¹. Other samples in the literature have lower carrier density and higher mobilities.⁵⁻¹⁶ These studies also show samples with higher carrier concentration ∼10¹⁹ cm⁻³; however, their bulk mobility is at least a factor of 2 higher. Apparently, all these reported mobilities are high enough to obscure surface oscillations. Moreover, our reference (nonstoichiometric) samples with n ∼ 10¹⁷ cm⁻³ have similar properties as other samples in the literature⁶ exhibiting 3D oscillations (see Appendix D).

The resistance versus magnetic field at 4.2 K is shown in Fig. 1. Strong oscillations are observed. Below we analyze the oscillations and show evidence that they arise from the topological surface states.

According to the Onsager relation

\[ F = \frac{h}{2\pi e} A(\epsilon_F), \]

the frequency of the SdH oscillations F is proportional to A(ε_F), the cross section of the Fermi surface in the plane perpendicular to the applied magnetic field. h is Planck's constant and e is the electron charge. For the case of a two-dimensional (2D) Fermi surface, \( F \propto \frac{1}{\sin \theta} \), where \( \theta \) is the angle between the normal to the 2D plane and the direction of the magnetic field. In Fig. 2 we show that our data follows the expected behavior of a 2D Fermi surface. Here \( \theta \) is the angle between the C₃ axis and the applied magnetic field. For \( \theta = 0 \), we find frequencies in the range of 190–198 T for S₁–S₅.

The behavior shown in Fig. 2 is in strong contrast with measurements of our low carrier-concentration samples. For these flakes, \( F = 20.26T \) for \( H \parallel C_3 \) and \( H \perp C_3 \), respectively. This corresponds to a 3D carrier density of 6.6 × 10¹⁷ cm⁻³, consistent with Hall measurements (data are presented in Appendix D).

Substituting the average frequency \( F = 195T \) in Eq. (1) and assuming a circular Fermi surface, we find \( k_F = 0.077 ± 0.003 \) Å⁻¹, which corresponds to a 2D carrier density of \( n_{2D} = \frac{k_F^2}{\pi} = 4.7 \times 10^{12} \) cm⁻², assuming no spin degeneracy, as should be the case for a topologically protected surface.

Figure 3 presents ARPES data from a cut taken along the Γ’M direction and going through the Γ point on a typical

---

**Fig. 1.** (Color online) Resistance as a function of field for a high carrier concentration sample (S1) at \( T = 4.2 \) K in a perpendicular magnetic field up to 32.5 T. Left inset: typical longitudinal resistance versus temperature at \( H = 0 \). Right inset: Hall resistance as a function of magnetic field at \( T = 2 \) K, \( H \parallel C_3 \) axis. The dashed line is a linear fit yielding a charge carrier concentration of 5.6 × 10¹⁹ cm⁻³. From the sign of the Hall slope we find that the charge carriers are electrons.

**Fig. 2.** (Color online) Frequency of the oscillations measured in S2 as a function of \( \theta \) (all other samples from this batch exhibited similar behavior). The bars represent the full width at half maximum of the FFTs, which are an upper bound for the error. The solid line is the expected 2D \( \frac{1}{\sin \theta} \) behavior.
The contribution of the bulk electrons as well as the Dirac-like surface state are clearly seen. We find \( k_F = 0.084 \pm 0.005 \text{Å}^{-1} \) in agreement with the value extracted from the SdH analysis. The resistance of a 2D system exhibiting SdH oscillations is given by \( R_s^{2D} = R_0 \left[ 1 + R_T \frac{\exp \left\{ 2\pi \left( \frac{F}{B} + \frac{1}{2} + \beta \right) \right\}}{\sinh \left( \frac{\alpha T}{B} \right)} \right] \), \( \text{(2)} \)

where \( R_0 \) is the zero field resistance and \( 2\pi \beta \) is the Berry phase which is expected to be \( \pi \) for an electron rotating around the Dirac point in a topological insulator at the low magnetic field limit. \( R_T \) contains the temperature dependence,

\[
R_T = \frac{\alpha T}{B} \left/ \sinh \left( \frac{\alpha T}{B} \right) \right.,
\text{(3)}
\]

with \( \alpha = \frac{2\pi \hbar \tau_D}{Me} \). The dingle factor is

\[
R_D = \exp \left( \frac{-\pi}{\omega_c \tau_D} \right),
\text{(4)}
\]

where \( \tau_D \) is the dingle scattering time, and corresponds to the dephasing of the Landau states. \( \omega_c = \frac{eB}{\hbar} \) is the cyclotron frequency. \( R_D \) determines the amplitude decay with the decrease of magnetic field.

The standard analysis of SdH oscillations yields a cyclotron mass of \( m^* \approx 0.16 m_e \) (Ref. 23; see Fig. 4); this together with the field dependence of the oscillation amplitude at 4.2 K gives \( \tau_D = (3.6 \pm 0.4) \times 10^{-14} \text{s} \), a Dingle temperature of \( T_D = \frac{\hbar}{2\pi \tau_D \hbar} = 33.78 \pm 0.33 \text{K} \), and a corresponding quantum mobility of \( \mu_{Q}^{2D} = 400 \text{cm}^2/\text{Vs} \).

Further confirmation of the 2D nature of the oscillations reported here is inferred from a comparison between the quantum (Dingle) and transport mobilities. Forcing a 3D fit to the amplitude of the oscillations (which contains an additional \( \sqrt{B} \) coefficient), one finds a 3D Dingle time of \( 3.9 \times 10^{-14} \text{s} \) which corresponds to a bulk quantum mobility of \( \mu_{Q}^{3D} = 570 \text{cm}^2/\text{Vs} \), larger than the transport mobility \( \mu_{T}^{3D} = 400 \text{cm}^2/\text{Vs} \). This is in strong contrast with recent reports for Bi\(_2\)Se\(_3\), where the Dingle mobility is four times smaller than the transport one.\(^{17}\) Furthermore, it is physically impossible to have \( \mu_{Q}^{3D} > \mu_{T}^{3D} \) in a bulk material, since the Dingle mobility takes into account all scattering events, including small angle scattering, that usually do not affect the resistivity, whereas for the transport mobility, backscattering events play the major role.

One of the hallmarks of a topological insulator is the \( \pi \) Berry phase associated with the Dirac dispersion relation. Since pronounced oscillations are observed in a broad field range (see Fig. 1), we are able to follow the evolution of the Berry phase as a function of applied magnetic field for S1.

The standard phase analysis is done by plotting \( F/B \) versus \( n \), the oscillation index (usually referred to as the Landau-level fan diagram). Applying this to our data, using the FFT peak \( F = 198 \text{T} \) yields a zero phase (\( \beta = 0 \)). We note that due to the many oscillations measured, the Fourier transform analysis favors a frequency that yields a zero phase. Furthermore, we find that \( \beta \) is very sensitive to \( F \), and even a deviation of 1.5%, which is well within the error margin, has a strong effect on \( \beta \). For example, if we take \( F = 195 \text{T} \), we get a \( \pi \) Berry phase (\( \beta = 0.5 \)) at low fields, deviating towards zero as the field increases. In the following, we study the behavior of \( \beta \) for frequencies in the vicinity of the FFT peak, \( F = 198 \pm 3 \text{T} \).

In order to focus on the field and frequency dependencies of the phase, we use the relation

\[
\beta(F,n) = n - F/B
\text{(5)}
\]

and plot \( \beta \) versus \( n \) in Fig. 5, where \( n \) is the index of the oscillation minima for the two frequencies mentioned (see Appendices A and B for more details and an analysis on more samples). The two behaviors mentioned above are now clearly seen. This analysis casts strong doubts on the
The work at Tel Aviv University is supported by the Israel Science Foundation (ISF) under Grant No. 1421/08 and the Ministry of Science and Technology. Part of this work has been supported by EuroMagNET under the EU Contract No. 228043.

APPENDIX A: ZEEMAN EFFECT ON $\beta(n)$

We use a simplified model of noninteracting electrons on the surface of topological insulators in a perpendicular external magnetic field. These electrons occupy orbital Landau levels which are coupled to their spin. The spin is also Zeeman coupled to the magnetic field. The energy of an electron in the $n$th orbital Landau level is therefore composed from both the orbital energy and the Zeeman energy in the following way (see, for example, Ref. 21):

$$E_n = \frac{1}{4} (g \mu_B B)^2 + 2\hbar^2 \left( \frac{v_F}{l_B} \right)^2,$$

(A1)

where $g$ is the effective gyromagnetic ratio for electrons on the surface, $\mu_B$ is the Bohr magneton, $B$ is the magnetic field, $v_F$ is the Fermi velocity associated with the Dirac cone, and $l_B = \sqrt{\frac{\hbar}{eB}}$ is the magnetic length. The $n$th maximum of the SdH oscillation in $R_{xx}$ occurs at magnetic field $B_{n,max}$ where this energy crosses the Fermi energy, $E_n = E_F$. This requirement and Eq. (A1) lead to

$$GB_{n,max} + n = \frac{F_0}{B_{n,max}},$$

(A2)

where $G = \frac{1}{2e}(\frac{L \mu_B}{v_F})^2$ and $F_0 = E_F^2/2hev_F^2$. One can write a similar expression for the minimum of the oscillations which occurs halfway between the $n$th and $n+1$ maxima. Interpolation of Eq. (A2) gives approximately (for large enough values

FIG. 5. (Color online) Analysis of the Berry phase $2\pi\beta$ according to Eq. (5) using the data from Fig. 1 after subtracting a smooth polynomial background. $n = 6.5$ corresponds to $B = 30.4$ T and $n = 28$ to $B = 7$ T. The blue diamonds are obtained using the FFT peak of $F = 198$ T, consistent with a zero Berry phase. The black circles correspond to $F = 195$ T, for which a $\pi$ Berry phase is observed at low fields. This phase changes towards zero as the field increases. This demonstrates the strong sensitivity of the phase to the frequency chosen. The error bars stem from the freedom in choosing the oscillation extrema. The green dashed (red dashed-dotted) line is the expected behavior of $\beta$ taking into account a Zeeman term with $g = 50$ (100), respectively.

IV. CONCLUSION

In summary, we present evidence that the surface states in Bi$_2$Se$_3$ can be probed in highly conducting flakes using the SdH effect. This is in contrast with the conventional approach focused on improving the crystal purity. From the SdH analysis we find for the surface states $m^* = 0.16m_e$, $\tau_D = (3.6 \pm 0.4) \times 10^{-14}$ s. The Fermi momentum in agreement with the ARPES data obtained on a flake from the same batch. We carefully study the behavior of the Berry phase with magnetic field, and show that two scenarios are possible: $\beta = 0.5$ ($\pi$ Berry phase) at low magnetic fields which changes fairly quickly with magnetic field, or a trivial zero phase. Two major issues are yet to be understood: first, the survival of protected surface states despite the large bulk conductivity, and second, the peculiar behavior of the Berry phase with magnetic field.

FIG. 6. (Color online) The behavior of $\beta(n)$ for S4 (green circles—196 T and squares—192.8 T) and S5 (blue ‘x’s—186.5 T and triangles—190 T).
of $n$),
\[ GB_{n_{\text{min}}} + n + \frac{1}{2} \simeq \frac{F_0}{B_{n_{\text{min}}}}. \]

The phase of the oscillations is defined as in Eq. (5),
\[ \beta(F,n) = n + \frac{1}{2} - \frac{F}{B_{n_{\text{max}}}} \simeq n - \frac{F}{B_{n_{\text{min}}}}. \tag{A3} \]

In usual metals, this definition gives $\beta = 0$, independent of $n$, for the right choice of $F$. However, in our case, given the theoretical relation between $n$ and $B_{n_{\text{max(min)}}}$ in Eqs. (A2) and (A3), $\beta$ depends on $n$ for any choice of $F$. We can choose $F = F_0$, which is the frequency of the oscillations in the limit of low magnetic fields, and we get for $n \gg 1$,
\[ \beta(F_0,n) = 0.5 + GB_{n_{\text{max}}} \simeq GB_{n_{\text{min}}}. \]

**APPENDIX B : FREQUENCY DEPENDENCE OF $\beta(n)$ FOR ADDITIONAL SAMPLES**

In Fig. 6 we show the frequency dependence of $\beta(n)$ for additional samples—S4 and S5 (measured up to 33 and 14 T, respectively). One can see that this is similar to the behavior of S1 depicted in Fig. 5.


dyodagan@post.tau.ac.il

1M. Z. Hasan and C. L. Kane, Rev. Mod. Phys. 82, 3045 (2010).


23. The cyclotron mass should not be zero even for the case of a linear dispersion relation. It is determined from the relation $m^* = E_F / v_F^2$ (Ref. 22).