The recent discoveries of surface and interface superconductivity with exceptionally high superconducting (SC) transition temperatures in several material structures [1–3] have drawn much attention to the phenomenon of strong type-II superconductivity in two-dimensional (2D) electron systems, in which the application of high magnetic fields can lead to exotic phenomena both in the normal and SC states [4]. Of special interest is the unique situation of the 2D superconductivity realized in helical surface states of topological insulators, e.g., Sb2Te3 [5], where the chemical potential $E_F$ is close to a Dirac point [6] (with Fermi velocity $v$) and the cyclotron effective mass $m^* = E_F / v^2$, [7] is a small fraction (e.g., 0.065 in Sb2Te3) of the free electron mass $m_e$ [8], resulting in a dramatic enhancement of the cyclotron frequency, $\omega_c = eH / m^* c$, and the corresponding Landau level (LL) energy spacing [9]. Furthermore, the spin-momentum locking, which is robust against non-magnetic impurity scattering [10], greatly enhances carrier mobility across the surface [5], resulting in relatively small LL broadening.

It should be noted, however, that the use of the standard LL spectrum, arising from a parabolic band-structure, in the self-consistent Bogoliubov-de Gennes theory, presented in ref. [9], has been done heuristically without actual derivation from the helical surface state Hamiltonian. Such a derivation is particularly necessary for the spin-momentum–locked model under study since SC pairing involves certain spin-orbital correlations. Our purpose in the present letter is therefore, two-fold: First, to develop the formal framework for solving the self-consistency equation for the SC order parameter in the 2D massless Dirac (Weyl) model Hamiltonian under a strong perpendicular magnetic field, and then exploit the developed formalism in a study of the transition to superconductivity in comparison with the well-known results of the standard model [4]. The SC transition in helical surface states of topological insulators, such as those reported, e.g., in ref. [5], is then comparatively studied with respect to both models. It is found that the calculated $H$-$T$ phase diagram for the Weyl model in the semiclassical limit (i.e., for LL filling factors $n_F > 1$) can be directly mapped onto that found for the standard model, having the same Fermi surface parameters $E_F$ and $v$, and a cyclotron effective mass equal to $m^* = E_F / 2v^2$. Significant deviations from the predicted mapping are found only for very small $E_F$, when the Landau-Level filling factors are smaller than unity, and $E_F$ shrinks below the cutoff energy.

We consider a free electrons gas on the surface (with coordinates $x$-$y$) of a topological insulator under a magnetic field $\mathbf{H} = (0, 0, H)$ (vector potential in the Landau gauge, $\mathbf{A} = (-Hy, 0, 0)$), characterized by $E_F$ and $v$. Expressing, unless it is otherwise explicitly defined, all length variables in units of the magnetic length $a_H \equiv \sqrt{eH / c}$, and all
energies in units of the cyclotron energy \(\hbar \omega_c \equiv \sqrt{2} \hbar v_F / a_H\), we adopt an effective mean-field Hamiltonian for singlet pairing in the Nambu representation \([10]\) (see the Supplemental Material (SM) Supplementary material.pdf, part A):

\[
\hat{H} = \left( \begin{array}{ccc} \sqrt{\gamma} \hat{\mathbf{\sigma}} \cdot \hat{\mathbf{p}} - \mu & i \hat{\mathbf{\sigma}} \Delta (\mathbf{r}) \\ -i \hat{\mathbf{\sigma}} \Delta^* (\mathbf{r}) & -\sqrt{\gamma} \hat{\mathbf{\sigma}} \cdot \hat{\mathbf{p}} + \mu \end{array} \right),
\]

where \(\mu \equiv E_F / \hbar \omega_c\), \(\hat{\mathbf{\sigma}} \equiv (\hat{\sigma}_x, \hat{\sigma}_y, \hat{\sigma}_z)\), with \(\hat{\sigma}_x, \hat{\sigma}_y, \text{ and } \hat{\sigma}_z\) the \(2 \times 2\) Pauli matrices, and unity matrix, respectively, \(\hat{\mathbf{p}} \equiv -i \nabla + (e/c) \mathbf{A}\) the gauge invariant momentum (charge -e), \(\Delta^* (\mathbf{r}) \equiv -|V| |\psi^0_\uparrow (\mathbf{r}) \psi^0_\downarrow (\mathbf{r})| \equiv \Delta^* (\mathbf{r}) = -\Delta^* (\mathbf{r})\) is the spin-singlet order parameter \([11]\), \(|\tilde{V}| \equiv |V| / \hbar \omega_c\), and \(V\) is the effective electron-electron interaction potential (given in units of energy \(\times\) area). Note that Zeeman spin-splitting is neglected with respect to the cyclotron energy in eq. (1) due to the very small cyclotron effective mass considered here.

The SC order parameter satisfies the self-consistency condition: \(\Delta^* (\mathbf{r}) = \tilde{V}(\mathbf{r}) \int d\mathbf{r}' \Delta^*(\mathbf{r}'|\mathbf{r}) Q(\mathbf{r}', \mathbf{r})\), where leading order in \(\Delta^* (\mathbf{r})\) the kernel is given by (see the SM, part C):

\[
Q(\mathbf{r}', \mathbf{r}) = \tau \sum_{\nu=\pm \infty} \left[ G^{(0)}(\mathbf{r}', \mathbf{r}; -\omega_e) G^{(0)}(\mathbf{r}', \mathbf{r}; \omega_e) 
- G^{(0)}(\mathbf{r}', \mathbf{r}; -\omega_e) G^{(0)}(\mathbf{r}', \mathbf{r}; \omega_e) \right],
\]

\(\tau \equiv k_B T / \hbar \omega_c\). \(\tau \equiv k_B T / \hbar \omega_c\). Note that the first product of propagators within the square brackets in eq. (2) corresponds to pairing of electrons that, due to spin orbit coupling, flip their spin projections, whereas the second product corresponds to pairing of electrons which preserve their spin projections. The normal-state Green’s functions, in the imaginary (Matsubara) frequency representation, with spin projections \(\alpha, \beta\), \(G^{(0)}(\mathbf{r}', \mathbf{r}; \omega)\) are elaborated in the SM, part B.

Similar to the situation in the standard, single-band (with quadratic energy dispersion) 2D electron system \([4,12]\), the order parameter of the ground Landau orbital: \(\Delta^* (\mathbf{r}) \sim e^{-2a q_F} / \sqrt{2} |y - q_x|\) is an eigenfunction of the integral operator with the kernel \(Q(\mathbf{r}', \mathbf{r})\), i.e.,

\[
\int d\mathbf{r}' Q(\mathbf{r}', \mathbf{r}) \Delta^*(\mathbf{r')} = A \Delta^*(\mathbf{r}),
\]

(see the SM, part E), so that the self-consistency equation can be reduced to the simple algebraic equation, 1 = \(|\tilde{V}| A\). Performing the summation over \(\omega_e\), we arrive at the following expression for \(|\tilde{V}| A\) (see the SM, part E):

\[
|\tilde{V}| A = \frac{1}{32} \frac{\lambda}{n_F} \sum_{ij=1}^{2} \sum_{N^{(i)}_n N^{(j)}_m} \frac{(m+n)!}{2^{m+n} n! m!} f^{(i)(j)}_{nm},
\]

\[
I^{(i)(j)}_{nm} = \frac{[(-1)^{i} \sqrt{m} + (-1)^{j} \sqrt{n}]}{2^{m+n} n! m!} \tan \left( \frac{\pi + (-1)^{i} \sqrt{m} + (-1)^{j} \sqrt{n}}{2 \mu + (-1)^{j} \sqrt{m} + (-1)^{i} \sqrt{n}} \right)
\times \tan \left( \frac{\pi + (-1)^{j} \sqrt{n} + (-1)^{i} \sqrt{m}}{2 \mu + (-1)^{i} \sqrt{m} + (-1)^{j} \sqrt{n}} \right),
\]

\(I^{(i)(j)}_{00} = 0\).

Fig. 1: (Colour online) Schematic illustration of the Weyl model bands structure for a positive chemical potential, smaller than the cutoff energy, showing a pair of Landau levels in both, conducting (CB) and valence (VB), subbands at the cutoff energy measured from the Fermi energy.

where: \(\lambda = \sqrt{n_F} |V| / \pi a_H^2 \hbar \omega_c = |V| N(E_F) = |V|[(m^* / 2\pi \hbar^2)^2], n_F \equiv \mu^2, N(E_F) = E_F / 2\pi (\hbar v)^2\) being the single-electron density of states per spin projection per unit area and \(m^* = E_F / v^2\) the effective cyclotron mass at the Fermi energy.

The different cutoff LL indices \(N^{(i)}_u (N^{(j)}_v)\), indicated in eqs. (3), refer to the different branches, i.e., the conduction (positive), or valence (negative) energy subbands of the Weyl model contributing to the pairing correlation. The different values arise due to the fact that the cutoff is introduced to the electron energy, by the mediating electron-phonon interaction, relative to the Fermi energy, rather than to the branching point (zero) energy of the Weyl bands structure. Thus, assuming a (Debye) cutoff energy \(\hbar \omega_D\), we should distinguish between two different situations. In the usual situation where \(\hbar \omega_D < E_F\), pairing takes place only in a single band, so that, e.g., for a positive chemical potential, we find: \(N^{(1)}_u = [n_F(1 + \gamma)^2], N^{(1)}_v = [n_F(1 - \gamma)^2]\), \(N^{(2)}_u = N^{(2)}_v = 0\), where \(\gamma = \hbar \omega_D / E_F\). In the unusual situation where the cutoff energy, \(\hbar \omega_D > E_F\) (see fig. 1), both inter and intra-band pairing take place, so that the cutoff LL indices are different for energies in the valence (VB) and conduction (CB) bands. Thus, for CB pairing (corresponding to the energy denominator \(2\mu - \sqrt{m} - \sqrt{n}\) in eq. (3)), we have: \(N^{(1)}_u = [n_F(1 + \gamma)^2], N^{(1)}_v = 0\). For the inter-band pairing (energy denominators \(2\mu - \sqrt{m} + \sqrt{n}\), or \(2\mu + \sqrt{m} - \sqrt{n}\) in eq. (3) the cutoff indices are: \(N^{(1)}_u = [n_F(1 + \gamma)^2], N^{(1)}_v = 0\), or: \(N^{(2)}_u = [n_F(1 - \gamma)^2], N^{(2)}_v = 0\), respectively.

In the semiclassical limit of our theory when the LL index at the Fermi energy, \(n_F\), is sufficiently large compared to unity, we may expand the CB energy appearing in the dominant contribution to \(A\) (i.e., \(\lambda^{(11)}_{nm}\)) in eq. (3) around \(m = n_F\), or \(n = n_F\), such that to first order in
(m - n_F)/2\sqrt{\tilde{n}_F} and (n - n_F)/2\sqrt{\tilde{n}_F}:

\[ |\tilde{V}|A \approx |\tilde{V}|A^{SC} = \frac{1}{8} \lambda S \sum_{m,n=N^{(1)}}^{N^{(1)}} \frac{(m+n)!}{2^{m+n+n!m!}} \times \frac{\tanh \left( \frac{n_F - m}{2\sqrt{\tilde{n}_F}} \right) + \tanh \left( \frac{n_F - n}{2\sqrt{\tilde{n}_F}} \right)}{2n_F - m - n}. \] (4)

Equation (4) is seen to be identical to the well-known expression for the pairing energy eigenvalue [4,13], provided that one makes the following replacements: \( \lambda \rightarrow 2\lambda_S, (2\sqrt{\tilde{n}_F}) \rightarrow \tau_S, N^{(1)} \rightarrow [n_F(1 + \gamma)], \) where \( \lambda_S \equiv |V|/(m_S^2/2\pi \hbar^2). \) This identity indicates that one can construct a useful reference for the 2D Weyl model developed above, in terms of a standard model, characterized by the (quadratic) energy-momentum dispersion, \( E = k^2/2m^*_c, \) and band effective mass, \( m^*_c, \) set equal to \( m^*_0/2 = E_{F0}/2\nu^2, \) with the same Fermi energy \( E_{F0} \) and wave number \( k_{F0} \) in both models, determined at a certain doping level in a concrete experiment (see below).

Under these assumptions: \( E_{F0} = \hbar^2k_{F0}^2/2m^*_0 = \hbar k_{F0}, \) and the standard cyclotron frequency, \( \omega^\infty_{c} \equiv eH/m^*_c, \) is related to the Weyl cyclotron frequency, \( \omega^W_{c} \equiv \sqrt{2\nu/aH}, \) via: \( \omega^W_{c} = 2\sqrt{\tilde{n}_F}(eH/m^*_0)^{1/2} = \sqrt{\tilde{n}_F}\omega^*_{c}, \) where in both models \( n_{F0} \equiv E_{F0}/\hbar\omega^\infty_{c} \equiv (E_{F0}/\hbar\omega^W_{c}) = (a_Hk_{F0})^2/2. \)

Using the set of parameters defined above, the well-known expression for the pairing energy eigenvalue obtained in the standard model takes the form:

\[ |\tilde{V}|A_S = \frac{1}{4} \lambda_S \sum_{m,n=n_{F0}(1-\gamma)}^{n_{F0}(1+\gamma)} \frac{(m+n)!}{2^{m+n+n!m!}} \times \frac{\tanh \left( \frac{n_F - m - 1/2}{2\tau_S} \right) + \tanh \left( \frac{n_F - m - 1/2}{2\tau_S} \right)}{2n_{F0} - n - m - 1}, \] (5)

where \( \tau_S = k_B T/\hbar\omega^\infty_{c}, \) and \( \gamma_0 = \hbar\omega_D/E_{F0}. \)

There is, however, an essential difference between the two models, and that is both \( \lambda \) and \( m^*_c \) in the Weyl model depend on the Fermi energy, whereas in the standard model both are constants. It will be, therefore, helpful to extend the reference model, expressed in eq. (5), for varying values of \( E_{F0}, \) by replacing \( n_{F0} \) with \( n_F \) as defined in the Weyl model, i.e., \( n_F \equiv (E_{F0}/\hbar\omega^W_{c})^2 = (a_Hk_{F0})^2/2, \) so that

\[ |\tilde{V}|A_S \rightarrow \frac{1}{4} \lambda_S \sum_{m,n=n_{F0}(1-\gamma)}^{n_{F0}(1+\gamma)} \frac{(m+n)!}{2^{m+n+n!m!}} \times \frac{\tanh \left( \frac{n_F - m - 1/2}{2\tau_S} \right) + \tanh \left( \frac{n_F - m - 1/2}{2\tau_S} \right)}{2n_{F0} - n - m - 1}. \] (6)

Experimental evidence for the existence of strong type-II superconductivity in a surface state of a topological insulator under a strong magnetic field can be found in results of transport, magnetic susceptibility, de Haas van
Fig. 3: (Colour online) $H$-$T$ phase diagram, obtained by solving the self-consistency equation for both models at the reference point: $\tilde{n}_F = \tilde{n}_{FO} = 10$, on the basis of the reference parameters $H_0$, $T_0$, as discussed in the text. The numerical scales of the abscissa and ordinate were obtained from the respective reduced variables $t$ and $h$. The cutoff was set at: $\hbar \omega_d/E_F = 0.5$. Deviations are seen only around the dark blue area, where the Weyl phase boundary is slightly above the standard one. Mutual reentrances of the SC and N phases, due to strong magneto-oscillations effect (see ref. [4]), are seen around the upper-left corner of the phase diagram.

Alphen (dHvA) oscillations and scanning tunnelling spectroscopy measurements, reported recently for Sb$_2$Te$_3$ [5]. Using a simple $s$-wave BCS model, similar to the standard model described above, with the experimentally observed dHvA frequency, $F_0 = 36.5$T (implying $n_{F0}(H) = F_0/H$), and cyclotron mass $m_s^* = 0.065m_e$, it was shown in [9] that such an unusual SC state can exist only in the strong-coupling superconductor limit. Here we study the relationships between the Weyl model and the extended standard model, described above, in the general parameters range, finding conditions for a complete mapping between the two models, and searching for physical situations in which they are qualitatively distinguishable. In fig. 2 we plot results of the pairing eigenvalue $A$, calculated within both models, as a function of the reduced magnetic field, $h \equiv H/H_0$, at the reduced temperature $t = T/T_0 = 0.01$, for various values of $E_F$. The temperature was selected sufficiently small to unfold the quantum oscillations associated with the Landau quantization.

Selecting for the reference parameters the values extracted from the transport and magneto-oscillations measurements [5]: $F_0 = 36.5$T, $H_0 = 2.5$T, $m_s^* = 0.065m_e$, and from the magnetic susceptibility measurements [5] the value: $T_0 = 100$ K, we define the dimensionless reference parameters: $\tilde{\tau}_S \equiv (k_B T_0/\hbar \omega_s^2)$ and $\tilde{\tau}_W \equiv (k_B T_0/\hbar \omega_W^2)$, where $\omega_s^2 \equiv (eH_0/m^*_S c)$ and $\omega_W^2 \equiv (\sqrt{2}c/a_H)$, so that $\tau_S = \tilde{\tau}_S(t/\hbar)$ and $\tau_W = \tilde{\tau}_W(t/\sqrt{\hbar})$. The two scales are therefore related via: $\tilde{\tau}_S = (\tilde{n}_{F0})^{1/2}\tilde{\tau}_W$, where $\tilde{n}_{F0} \equiv (a_H k_F)^2/2 \approx 10$.

The eigenvalues $A_W$, $A_S$, plotted in fig. 2 as functions of $h$, for various values of $\tilde{n}_F \equiv (a_H k_F)^2/2$, show at $\tilde{n}_F = \tilde{n}_{FO}$ complete agreement between the two models, including the fine structure of the quantum oscillations. Under these conditions, solutions of the self-consistency equation, $1 = |\tilde{V}|A$, in both models, yield nearly identical results for the $H$-$T$ phase diagrams, as shown in fig. 3, except for a small deviation in the low fields region, due to the different ultraviolet divergency predicted by the two models. The two intersection points of the phase boundary with the axes, shown in fig. 3, are seen to be close to $t = 1$ and $h = 1$, thus indicating that the calculated $H_{c2}(T \rightarrow 0)$ and $T_c(H \rightarrow 0)$ values are close to the values of $H_0$ and $T_0$, respectively.

For values of $\tilde{n}_F$ away from $\tilde{n}_{FO}$ the baseline of $A_W$ is shifted with respect to that of $A_S$, depending on whether $\tilde{n}_F > \tilde{n}_{FO}$ (shift up), or $\tilde{n}_F < \tilde{n}_{FO}$ (shift down), thus reflecting the dependence of the pairing correlation in the Weyl model on the carrier density through the dependence of its single-electron density of states on the Fermi energy. Indeed, comparing the semiclassical expression, eq. (4), to the standard eigenvalue, $A_S$, given in eq. (6), one finds for doping levels away from the reference point, i.e., for $k_F \neq k_{FO}$, the simple relation: $A_W^{SC} = (k_F/k_{FO})A_S$.

Note, however, that the oscillatory patterns of $A_W$, $A_S$ in fig. 2 remain nearly the same, except for slight relative narrowing of the Weyl peaks upon decreasing $\tilde{n}_F$, which becomes quite significant in the quantum limit, e.g., at $\tilde{n}_F = 0.5$ in fig. 2(d).

It is also remarkable that in the ultimate quantum limit, i.e., when $\tilde{n}_F \rightarrow 0$, the pairing correlation in the Weyl model, despite its vanishing normal electron density of states at the Fermi energy, does not vanish (see the inset in fig. 2(d)). This non-vanishing limiting behavior, which is an apparent outcome of the fact that the prefactor $\lambda/\sqrt{\pi} = |V|/\pi a_H^2 \hbar c$, in eq. (3) is independent of the Fermi energy, is similar to that found for the standard model. The much smaller limiting value obtained for the Weyl model is due to the absence of the zero-zero Landau level term $I_{00}^{(ij)}$ in eq. (3). Note that this peculiar behavior is inherent to the mean-field approximation in the 2D model systems under study. Thus, critical SC phase fluctuations in the limit of zero carrier density [14] are expected to destroy long-range SC order in both models.

To summarize, we have developed a Nambu-Gorkov Green’s function approach to strongly type-II superconductivity in a 2D spin-momentum–locked (Weyl) Fermi gas model at high perpendicular magnetic fields in order to study the transition to high-field surface superconductivity observed recently on the topological insulator Sb$_2$Te$_3$ [5]. We have found that, for LL filling factors larger than unity, superconductivity in such a 2D Weyl Fermion gas can be mapped onto the standard 2D electron (or hole) gas model, having the same Fermi surface parameters, but with a cyclotron effective mass, $m^* = E_F/2v^2$, which could be dramatically reduced below the free electron mass, $m_e$, by manipulating the doping level, or the gate voltage. Our calculations for Sb$_2$Te$_3$ show that the SC helical surface state reported in [5] was in the moderate semiclassical range ($\tilde{n}_F \gtrsim 10$), so justifying the mapping with the standard model. They reveal a very unusual, strong type-II superconductivity at low carrier density and
small cyclotron effective mass, $m^* = 0.065m_e$, which can be realized only in the strong-coupling ($\lambda \sim 1$) superconductor limit [9]. Further reduction of the carrier density in such a system could yield an effective cyclotron energy comparable to or larger than the Fermi energy, LL filling factors smaller than unity, and cutoff energy larger than the chemical potential, resulting in significant deviations from the predictions of the standard model.

However, for such a dilute ferromagnet gas system the simple mean-field BCS theoretical framework of superconductivity, exploited in this paper, should be drastically revised, particularly due to the neglect of both phase and amplitude fluctuations of the SC order parameter [14], and to the breakdown of the adiabatic approximation in the electron phonon system [15]. Several recent reports on superconductivity in very dilute ferromagnet gas systems, such as that found in compensated semimetallic FeSe [16], or in the large-gap semiconductor SrTiO$_3$ [17], have drawn much attention to fluctuation superconductivity beyond the Gaussian approximation, which could lead to crossover between weak-coupling BCS and strong-coupling Bose-Einstein condensate limits [18]. In the presence of strong magnetic fields the situation is further complicated due to complex interplay between vortex and SC amplitude fluctuations [19].

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