Fermionic analogue of black hole radiation with a super high Hawking temperature

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Abstract: Measurement of gravitational Hawking radiation of black hole (BH) is prohibitive because of an extremely low Hawking temperature ($T_H$). Efforts have been devoted to creating artificial BHs with a higher $T_H$ but having rather limited success. Here we demonstrate a fermionic analog of BH with a super high $T_H \sim 3$ K. We show that Floquet-Dirac states, formed in a periodically laser driven black phosphorous thin film, can be designed with a spatial gradient to mimic the “gravity” felt by fermionic quasiparticles as that for a Schwarzschild BH (SBH). Quantum tunneling of electrons from a type-II Dirac cone (inside BH) to a type-I Dirac cone (outside) emits a SBH-like Hawking radiation spectrum. This work provides a laboratory design of fermionic BH and a condensed-matter analogue to study fascinating astrophysical phenomena.
Gravitational black hole (BH) is a curved spacetime absorbing everything inside its event horizon because of an extremely large mass/radius ratio, as characterized by overtilted light cones according to general relativity theory. This notion of nothing escapes from a BH is no longer valid if quantum effects are considered. According to Hawking [1-3], quantum fluctuations at the event horizon generate particles and antiparticles to propagate out of and into the BH, respectively. Consequently, a BH evaporates thermally like a black body, known as Hawking radiation, as shown in the upper panel of Fig. 1. The intensity of this quantum radiation is quantified by Hawking temperature \( T_H = \kappa / 2\pi \) with the gravity \( \kappa \) at the event horizon, which is however extremely low, \( \sim 10^{-8} \) K for a BH with one solar mass [1]. This prohibitively weak intensity has prevented us from revealing the fundamental connection between quantum mechanics and gravitation. To overcome this difficulty, artificial analogs have been proposed including sonic BHs [4,5] in Bose-Einstein condensates [6-9], ion rings [10], and Fermi-degenerate liquids [11]. Unfortunately, \( T_H \) for most artificial BHs is still very low, e.g. \( \sim 10^{-9} \) K in Bose-Einstein condensates [7]. Furthermore, the observation for optical BH analogs in ultrashort laser pulse filaments remains controversial [12-21]. Therefore, new BH analogs inherent with quantum effect and a much higher \( T_H \) are highly desirable to detect Hawking radiation.

Recently, type-II Weyl/Dirac fermions in solids, with a unidirectional Fermi velocity, have been proposed as a new platform to realize artificial BHs [22-27]. In the regions where type-II and type-I fermions are separated by a boundary with type-III fermions (i.e. the “event horizon”), two worlds inside and outside a BH are analogously formed (lower panel in Fig. 1). Quantum mechanics enables Hawking radiation of quasiparticles at the event horizon. Thanks to the fact that electrostatic interactions in solids are orders-of-magnitude stronger than gravitational forces, an unprecedented high \( T_H \) is expected. Our previous work shows that laser-driven black phosphorus (BP) can host type-I, -II and -III fermions [24], manifesting a potential fermionic analogue of BH. In order to support artificial Hawking radiation in this real material, however, two requirements are necessary: (i) a spatial distribution of band structure to produce a “steep” analogous gravity field, and (ii) a working mechanism to induce fermionic Hawking radiation.
In this Letter, we propose an experimentally accessible solid-state fermionic analogue of BH in two-dimensional (2D) BP under laser illumination. Combining first-principles and quantum tunneling calculations, a spatially inhomogeneous system with successively distributed type-II, -III and -I Dirac fermions is designed, which acts like an analogous Schwarzschild BH (SBH) metric to induce electron emission from type-II to type-I region. An effective gravity field corresponding to a striking $T_{\text{H}} \sim 3$ K is achieved.

Dirac fermions are classified into different types, and successive transitions between them in 2D can be described by the Hamiltonian

$$H(k) = c_{s}k_{s}\sigma_{s} + c_{y}k_{y}\sigma_{y} + vk_{0}\sigma_{0},$$  \hspace{1cm} (1)

where $\sigma_{s} = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}$, $\sigma_{y} = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}$, and $\sigma_{0}$ is the identity matrix. Dispersions in $k_{y}$ direction are $\varepsilon_{1} = (v + c_{y})k_{y}$ and $\varepsilon_{2} = (v - c_{y})k_{y}$, whose crossing forms massless Dirac cone with Fermi velocities $v_{F1} = v + c_{y}$ and $v_{F2} = v - c_{y}$, respectively. $v_{F1}$ and $v_{F2}$ can be tuned by changing $v$ and $c_{y}$, leading to three types of cones with distinct band dispersions and Fermi surfaces. For an upright type-I Dirac cone, $|c_{y}| >> |v|$, $v_{F1} = -v_{F2}$. When the upright cone tilts, $|c_{y}|$ and $|v|$ decreases, and we elaborate here on the case of a clockwise tilt having $c_{y} > 0$ and $v < 0$, as shown in Fig. 2(a). The cone remains as type-I when $c_{y}$ decreases from $+\infty$ to $-v$ ($c_{y} > -v$).

When it is tilted to a critical point ($c_{y} = -v$), the cone has a flat band of $\varepsilon_{1}$ ($v_{F1} = 0$, $v_{F2} < 0$), which is dubbed type-III with a line-like Fermi surface [23,24]. Beyond the critical point, an overtilded cone is named type-II having $c_{y} < -v$, whose two Fermi velocities have the same sign ($v_{F1}$, $v_{F2} < 0$) and Fermi surface encompasses both electron and hole pockets.

We next transform the Dirac fermions in crystals to the particles in gravity field. In Einstein notation, Eq. (1) can be rewritten as $H(k) = e_{j}^{i}\sigma^{j}k_{i} + e_{0}^{i}\sigma^{0}k_{i}$. $i, j = 1, 2$ (or $x, y$) and the matrix $e_{j}^{i}$ and vector $e_{0}^{i}$ are equivalent to components of a tetrad field $e_{\alpha}^{\mu}$ ($\alpha, \mu = 0, 1, 2$ in $2 + 1$ dimensions) in general relativity. Then, an effective relativistic covariant metric $g_{\mu\nu} = (\eta^{\alpha\beta}e_{\alpha}^{\mu}e_{\beta}^{\nu})^{-1}$ with $\eta^{\alpha\beta} = \text{diag}(-1,1,1)$ governs Dirac fermions. The corresponding line element ($ds^{2} = g_{\mu\nu}dx^{\mu}dx^{\nu}$) is

$$ds^{2} = \left(1 - \frac{v^{2}}{c_{s}^{2}}\right)dt^{2} + \frac{1}{c_{s}^{2}}dx^{2} + \frac{1}{c_{s}^{2}}dy^{2} - \frac{2v}{c_{s}^{2}}dtdy,$$  \hspace{1cm} (2)
which shows the behavior of Dirac quasiparticles in an effective 2 + 1 dimensional \((t, x, y)\) spacetime. For \(v = 0\), \(ds^2\) in \((t, y)\) spacetime becomes \(-c_y^2 dt^2 + dy^2\), representing a flat spacetime where electronic wave propagates at the velocity of \(c_y\) along \(\pm y\) directions. Hence, \(c_y\) corresponds to light velocity \((c)\) in the gravitational spacetime, establishing that Dirac wave propagating in a crystal field is formally equivalent to light propagating in a gravity field. In contrast to a constant \(c\), however, the “light velocity” \((c_y)\) of Dirac fermions can be changed by interactions in a crystal.

Effective spacetime corresponding to a given type of Dirac fermion can be designed by the relative magnitude of \(c_y\) and \(v\). The Dirac cone manifests in its “spacetime” as an artificial light cone \((ds^2 = 0)\) with dispersions of \(t_1 = \frac{y}{v + c_y}\) and \(t_2 = \frac{y}{v - c_y}\), as shown in Fig. 2(b). For type-I, \(t_1' = 1/(v + c_y) > 0\) and \(t_2' = 1/(v - c_y) < 0\) having the opposite sign, so that quasiparticles propagate along both \(+y\) and \(-y\) directions. For type-III, \(t_1 \to \infty\), so that the one branch of quasiparticles stays at a fixed location (the event horizon) permanently; while the other branch propagates along \(-y\) direction \((t_2' < 0)\). For type-II, the minus sign of both \(t_1'\) and \(t_2'\) indicates that all quasiparticles propagate along \(-y\) direction, resembling the unidirectional behavior of particles inside a SBH.

To produce a desired gravity field to be felt by Dirac fermions, spatial distribution of effective geometry should be designed appropriately. Matching the effective and Schwarzschild metrics [28], the effective potential energy of quasiparticles is \(\phi(y) = -\frac{1}{2} \frac{v^2(y)}{c^2_y(y)}\). \(\Phi(y)\) is inversely proportional to the distance: \(\phi(y) = -\frac{1}{2} \frac{y_h}{y}\), where \(y_h\) is the location of event horizon. As a result, the effective light velocity \(c_y\) distributes along the \(y\) direction as

\[
c_y(y) = -v \sqrt{\frac{y}{y_h}}. \tag{3}
\]

This guarantees that type-II fermions in the \(0 < y < y_h\) region represents inside of a BH, type-III at \(y = y_h\) represents the event horizon, and type-I in the \(y > y_h\) region represents the outside of a BH [Fig. 2(c)]. The motion of quasiparticles can be described by kinematic equations of

\[
k_{y_1} = \frac{\epsilon_1}{v(1 - \sqrt{\frac{y}{y_h}})} \quad \text{and} \quad k_{y_2} = \frac{\epsilon_2}{v(1 + \sqrt{\frac{y}{y_h}})}.
\]

As shown in Fig. 2(d), \(k_{y_1}\) and \(k_{y_2}\) represent two ingoing
waves in the $0 < y < y_h$ region, and one ingoing and one outgoing wave in the $y > y_h$ region. At $y = y_h$, there is such a high potential barrier that quasiparticles occupying on $\varepsilon_1$ band are impossible to go through. This scenario in the classical perspective of general relativity, however, would break down because of the quantum fluctuations at event horizon, producing Hawking radiation.

Next, using quantum tunneling method \cite{29,30}, we analyze the analogous Hawking radiation at the effective BH event horizon with the curved geometry of Eq. (3). As illustrated in Fig. 2(c), inside the BH the $\varepsilon_1$ states above the Dirac point ($\varepsilon_1 > \varepsilon_D$) are excited, while outside the BH the $\varepsilon_1$ states are empty. This leads to emission of the excited electrons from inside to outside the BH. Adopting Wentzel-Kramers-Brillouin approximation, the tunneling probability is $P = 1/(1+\exp(2S))$ with a classical action $S = \text{Im} \int k_y(y)dy$. Under the approximation of $k_y \approx \varepsilon \left( \frac{dc}{dy} \right)_{y_h} \cdot (y - y_h)$ for the $k_y$ branch around the event horizon, $P$ produces a spectrum with the energy intensity $I(\varepsilon) = n(\varepsilon) \cdot \varepsilon$. The number of radiated electrons follows

$$n(\varepsilon) \propto \frac{\varepsilon^2}{\exp\left(\frac{\varepsilon}{k_BT_H}\right) + 1},$$

where $k_B$ is the Boltzmann constant, and $T_H = \frac{1}{2\pi k_B} \left| \frac{d(c,-\nu)}{dy} \right|_{y_h}$ is the Hawking temperature.

The spectrum $I(\varepsilon)$ conforms to the thermal radiation of a black body by Planck’s law. Namely, this spectrum of analogous Hawking radiation is the same as that of gravitational counterpart. In analogy with particles and antiparticles produced from gravitational Hawking radiation, the "radiated" electrons and holes created in a quantum fluctuation are entangled (lower panel in Fig. 1).

The predicted fermionic analogue of Hawking radiation can be measured by modern experiments, which should be much easier than previously proposed analogs \cite{6-8,10,11} and gravitational BHs. The Planck spectrum $I(\varepsilon)$ can be obtained by measuring the local distribution of electronic states $n(\varepsilon)$ in the region $y > y_h$, using scanning tunneling spectroscopy or angle resolved photoemission spectroscopy (ARPES). As shown in Fig. 2(e),
the spectrum \( n(\epsilon) \) with \( \epsilon > \epsilon_D \) has a peak at \( \epsilon_p = 0.6 \) meV for \( T_H = 3 \) K, which is distinctively different from the Fermi-Dirac distribution; while \( n(\epsilon) \) for \( \epsilon < \epsilon_D \) remains the same as before. The Planck spectrum in the region of \( \epsilon > \epsilon_D \) provides a key signature for detecting Hawking radiation. \( T_H \) can be tuned by controlling the effective gravity field to facilitate the detection of \( n(\epsilon) \). The relation between \( \epsilon_p \) and \( T_H \) follows 
\[
\frac{1}{2} \frac{\epsilon_p}{k_B T_H} = 1 + \exp \left( - \frac{\epsilon_p}{k_B T_H} \right),
\]
leading to a linear dependence \( \epsilon_p = 2.2 \cdot k_B T_H \) [Fig. 2(f)]. An unprecedentedly high \( T_H \) is achieved because interactions in crystals are orders-of-magnitude stronger than gravity field.

Based on above analysis, we next conceive a solid-state material to realize such fermionic analogue of BH. By applying a gate with the vertical electric field \( E_{ext} \) or compressive strain \( \delta \) along armchair (x) direction, the direct band gap of 2D BP decreases \([24,31-37]\), leading to an inversion \( \Delta \epsilon \) of valence \( \epsilon_1 \) and conduction \( \epsilon_2 \) bands. Symmetry-protected type-I Dirac cone emerges, which was confirmed by ARPES \([34,35]\). For bilayer BP, this Dirac state is shown in Fig. 3(c,d), where \( \Delta \epsilon = 14 \) meV at \( \Gamma \) point is induced by \( E_{ext} = 0.18 \) V/Å or \( \delta = 7.6\% \) (see Fig. S1 for details \([38]\)).

To form the three types of Dirac states by laser-driving, we study coherent interactions between bilayer BP and linearly polarized laser (LPL) with a time-dependent vector potential \( A(t) = A_0 \sin(\omega t, 0, 0) \) [Fig. 3(a)]. The photon energy of the time-periodic and space-homogeneous LPL is chosen as \( \hbar \omega = 0.03 \) eV, which is larger than \( \Delta \epsilon = 14 \) meV to avoid the crossing, nearby Dirac point, of the original \( n = 0 \) and photon-dressed \( n \neq 0 \) bands. When the LPL is applied, the type-I Dirac cone tilts due to the hybridization between \( n = 0 \) and \( n \neq 0 \) bands. At the critical laser amplitude \( A_0 = 16 \) V/c (corresponding to 0.24 mV/Å or \( 7.65 \times 10^5 \) W/cm², here \( c \) is light velocity), \( \epsilon_1 \) band along \( \Gamma Y \) path becomes flat, forming type-III Dirac cone [Fig. 3(e,f)]. As the laser amplitude increases to \( A_0 = 20 \) V/c, the Dirac cone tilts further to become type-II [Fig. 3(g,h)]. Then not only the slope of \( \epsilon_1 \) and \( \epsilon_2 \) dispersions has the same sign, but the states on \( \epsilon_1 \) band above Fermi level \( \epsilon_1 > 0 \) eV is also occupied by electrons, which agrees with the excited type-II cone in Fig. 2(c). Consequently, type-I, II and -III Dirac fermions are created in a single material of 2D BP, depending sensitively on laser intensity. In addition, laser frequency also plays a crucial role in determining the type of these cones,
which is shown in the phase diagram of Fig. 3(b).

To mimic the spacetime geometry of Eq. (2), we map the photoinduced Dirac states from \textit{ab initio} calculations to model parameters in Eq. (1). As shown in Fig. 3(i), the parameter \( \nu = -0.1 \text{ eV} \cdot \text{Å} \) is constant, while \( c_y \) decreases gradually with increasing \( A_0 \), showing the transition from type-I \((c_y > -\nu)\) to type-II \((c_y < -\nu)\). In the region of strong (weak) laser intensity, \( c_y \) and laser amplitude \( A_0 \) exhibit a linear (nonlinear) relation, which is a typical characteristic of optical Stark effect (see details in Fig. S3 [38]). Neglecting the minor deviation in the weak intensity regime, the relationship between \( c_y \) and \( A_0 \) can be fitted linearly by

\[
    c_y[\text{eV} \cdot \text{Å}] = 0.314 - 0.013A_0[V/c],
\]

(5)

which is the base to explore Hawking radiation in laser-driven bilayer BP.

Substituting \( c_y \) in Eq. (3) with Eq. (5), the amplitude of space-inhomogeneous laser field is

\[
    A_h(y)[V/c] = -8.154 \sqrt{\frac{y[A]}{y_s[A]} + 24.1538},
\]

leading to \( T_H = \frac{-\nu}{4\pi k_B T_H} \). To achieve \( T_H = 3 \text{ K} \), the BH size should be set at \( y_h = 30.2 \text{ Å} \), which requires the laser field to decrease from 0.37 to 0 mV/Å in a range of 260 Å along the zigzag \((y)\) direction [Fig. 4(a)]. This gradient of laser intensity can be readily realized in experiments [44-46]. As shown in Fig. 4(b), the peak position \( \epsilon_p \) of the experimental Hawking spectrum, \( I_{\text{exp}}(\epsilon) = n(\epsilon) \cdot \text{DOS}(\epsilon) \), reaches 0.6 meV above the Fermi level. It is clearly different from the spectrum of thermal excitation at \( T_t = 3 \text{ K} \) in Fig. 4(c). Decreasing the BH size, \( \epsilon_p \) increases in a wide range of \( T_H \) [Fig. 4(d)]. Also, by engineering Dirac states via changing photon energy, Hawking radiation can also be controlled. Finally, we note that there are a pair of Dirac cones along \(-Y\) path resembling black and white holes respectively, but their coexistence does not influence the Hawking radiation (see details in Fig. S4 [38]).

In conclusion, type-I, -II and -III Dirac fermions are predicted to form in BP thin film under LPL. When the laser field is spatially inhomogeneous, two regions with type-I and -II Dirac fermions are separated by the boundary with type-III Dirac fermions, resulting in the fermionic analogue of BH Hawking radiation. Most importantly, Hawking temperature can reach 3 K, with the emitted electrons occupying experimentally observable states peaked at 0.6 meV above the Fermi level. The Hawking radiation can be well controlled by laser field.
Our finding enables intentional engineering of fermions to provide a table-top condensed matter platform for simulating exotic phenomena in astrophysics and general relativity.

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FIG. 1. Schematic illustration of gravitational Hawking radiation and fermionic analog. Upper panel: Light cones indicate that particles cannot escape from gravitational BH classically (arrowed straight line), but particles and antiparticles can emit from event horizon due to quantum fluctuation (arrowed wavy line). Lower panel: Dirac cones show that electrons cannot escape from the type-II region (arrowed straight line) of the artificial BH in crystals, quantum tunneling enables emission of electrons and holes from the event horizon (type-III region) to type-I and -II regions respectively (arrowed wavy line).
FIG. 2. Fermionic analogue of a BH and consequent Hawking radiation. (a) Fermi velocity $v_F$ of Dirac cone varies with the parameter $c_y$, resulting in the transition from type-I (green), type-III (red) to type-II (blue) Dirac cone successively. (b) Three types of artificial light cone (green, red, and blue) formed via counterclockwise rotation, corresponding to type- I, -III, and -II Dirac fermions respectively. (c) The distribution of $c_y$ of type-I, -III, and -II Dirac fermions appears successively along $+y$ direction, producing the same “gravity field” felt by electrons as that of particles in a SBH. (d) Kinematical equation of motion shows a high potential barrier at $y_h$ to prevent electrons from escaping from type-II to type-I region, but this process takes place by quantum tunneling indicated by the arrowed red dashed line in (c) and (d). (e) The spectrum of electron occupation probability ($n$) produced by Hawking radiation (red line) vs. thermal excitation (gray line) at 3 K. (f) The relation between the peak position $\epsilon_p$ in (e) and Hawking temperature $T_H$. 
FIG. 3. Photoinduced type-I, -II and -III Dirac fermions in bilayer BP. (a) Laser field $A(t) = A_0(\sin(\omega t), 0, 0)$ is polarized along armchair ($x$) direction. (b) Dirac states induced by laser with different amplitude $A_0$ and frequency $\omega$. (c,d) Band structure of bilayer BP under vertical static electric field of 0.17 V/Å (or 7.6% compressive strain). Floquet-Bloch band structure of bilayer BP driven by laser with amplitude $A_0 = 16$ V/c (e,f) and $A_0 = 20$ V/c (g,h), and photon energy $\hbar\omega = 0.03$ eV. (i) With increasing laser intensity, type-I Dirac fermion is transitioned to be type- III, and -II successively. The gray dashed line shows linear fitting of $c_y$ vs. $A_0$. 

FIG. 4. Experimentally detectable and tunable Hawking radiation in bilayer BP. (a) Required laser intensity distribution along zigzag ($y$) direction to realize a SBH with $T_H = 3$ K. (b) The experimental spectrum $I_{\text{exp}}(\varepsilon) = n(\varepsilon) \cdot \text{DOS}(\varepsilon)$ of Hawking radiation with $T_H = 3$ K, which is compared with that of thermal excitation at $T_t = 3$ K (c). The spectrum in ground state is shown by gray dashed line. (d) The dependence of $T_H$ on the location of event horizon. The red dot corresponds to $T_H = 3$ K and $\varepsilon_p = 0.6$ meV.