A signature property of a large class of magnetized and unmagnetized plasmas in the Laboratory, Space, and Astrophysical systems is the extremely low collisionality that can be due to high plasma temperature $T_e$ or low plasma density $n_e$, or a combination of the two [1]. For example, a fusion-grade plasma in a tokamak reactor has $T_e \sim 10 - 20$ kilo-electron-volts (KeV) and $n_e \sim 10^{19-20}$ per cubic meter, which result in a mean-free-path $\lambda_{mf,p} \sim 10^4$ meters (m), while the toroidal length of the confinement chamber is merely 20-30 m [2–4]. In the earth’s radiation belt, the electron $\lambda_{mf,p}$ can be as long as $10^{11}$ m with electron energy from tens of KeV to MeV and an $n_e \sim 10^4$ m$^{-3}$ [5–8]. At the even grander scale of clusters of galaxies, the intracluster hot gas has $n_e \sim 10^5 - 10^4$ m$^{-3}$ and $T_e \sim 2 \times 10^5 - 10^6$ K [9–13], so $\lambda_{mf,p}$ is in the order of tens of kilo-parsec to mega-parsec.

A whole class of problems arises if a localized cooling spot is introduced into such a nearly collisionless plasma. This could be structure formation in a galaxy cluster where a radiative cooling spot is driven by increased particle density [9] or an event horizon of a black hole that provides an absorbing boundary for plasmas [14]. A satellite traversing the earth’s radiation belt can be a sink for plasma energy and particles [15–16]. In a tokamak reactor, solid pellets that are injected into the fusion plasma for fueling and disruption mitigation [17–19], provide localized cooling due to a combination of energy spent on phase transition and ionization of the solid materials, and the radiative cooling that is especially strong when high-Z impurities are embedded in the frozen pellet. Even in the absence of pellet injection, large-scale magnetohydrodynamic instabilities can turn nested flux surfaces into globally stochastic field lines that connect fusion core plasma directly onto the divertor/first wall [20–24], causing a thermal collapse via fast parallel transport along the field lines within a short period of time that can range from micro-seconds to milliseconds [24]. An outstanding physics question is how a thermal collapse of the surrounding plasma, commonly known as a thermal quench in tokamak fusion, would come about in such a diverse range of applications.

The most obvious route for the thermal collapse is via electron thermal conduction along the magnetic field line that intercepts the cooling spot, for which the Braginskii formula [25] would produce an enormous heat flux [9, 11] if there is a sizeable temperature difference $\Delta T = T_{e0} - T_{en} \approx T_{e0}$ between the cooling spot ($T_{en}$) and the surrounding plasma ($T_{e0}$),

$$q_{el} = n_e k_e \delta^2 \nabla T_e \sim n_e v_{th,e} \frac{\lambda_{mf,p}}{L_T} \Delta T \sim n_e v_{th,e} \frac{\lambda_{mf,p}}{L_T} T_{e0}. \tag{1}$$

Here $v_{th,e} = \sqrt{T_{e0}/m_e}$ is the electron thermal speed and $L_T$ the distance or field line length over which the temperature drop $\Delta T$ is established. For a nearly collisionless plasma, the temperature collapse necessarily starts with Knudsen number $K_n \equiv \lambda_{mf,p}/L_T \gg 1$, a regime in which the free-streaming limit [26] of

$$q_{el} \approx \alpha n_e v_{th,e} T_{e0}, \tag{2}$$

is supposed to apply in lieu of Braginskii, with $\alpha \approx 0.1$ [27]. The pressure-gradient-driven plasma flow $V_{\parallel 0}$ along the magnetic field line is limited by the ion sound speed $c_s$, so the convective electron energy flux is bounded by $n_e c_s T_{e0}$. The scaling of $q_{el}$ with $v_{th,e}$ in Eq. (2) suggests that the electron energy flux would be dominated by conduction as normally $v_{th,e} \gg c_s$ in a plasma of comparable electron and ion temperatures. In such a conduction-dominated situation, the much colder but denser cooling spot would be rapidly heated up by the electron thermal conduction from the surrounding hot plasma, and as the result, it can become over-pressured and the original cooling spot, say an ablated pellet in a tokamak, tends to expand into the surrounding plasma, yielding an outflow.

In the aforementioned problem of clusters of galaxies, one has instead observed robust cooling flows into the radiative cooling spot that aggregate mass onto the cooling spot [9], although more recent observations reveal a more modest mass-accreting cooling flow that indicates the role of various additional heating mechanisms to balance the cooling [11–13]. This is inconsistent with the conduction-dominated scenario.
which would overwhelm the convective energy transport (scaling of Eq. (4)). As the result, a robust constraint in a nearly collisionless plasma is realized by ambipolar transport. In the case of precooling, the boundary of which recycles all mass being entrained by the cooling flow apparent to aggregate mass towards the cooling spot. There are totally four fronts for the propagation of the electron fronts that originate from the cooling spot with characteristic speeds. The necessary constraint is on the spatial gradient of the electron parallel conduction flux, which can be seen from the energy equation for the electrons along the magnetic field,

\[ n_e \left( \frac{\partial}{\partial t} T_e + V_e \frac{\partial}{\partial x} T_e \right) + 2n_e T_e \frac{\partial}{\partial x} V_e + \frac{\partial}{\partial x} q_{en} = 0. \]  

(3)

Here \( x \) is the distance along the magnetic field line, \( n_e, T_e, V_e \) are the density, parallel temperature, and parallel flow of the electrons, and \( q_{en} \equiv \int n_e (v_e - V_e) d^3 \) is a component of the parallel heat flux. Let the cooling flow span a length \( L_T \), one can see the convective energy transport terms follow the scaling of \( n_e T_e V_e / L_T \). Ambipolar transport constrains \( V_e \approx V_{th} \approx m_i^{-1/2} \), so the free-streaming scaling of \( q_{en} \) in Eq. (2) would predict \( \partial q_{en} / \partial x \sim \alpha n_e v_{th} T_e / L_T \ll m_i^{-1/2} \), which would overwhelm the convective energy transport (\( \propto m_i^{-1/2} \)) to force a \( T_e \) collapse and remove the pressure gradient drive that sustains the cooling flow. The condition for accessing the cooling flow regime of plasma thermal quench is thus

\[ \frac{\partial q_{en}}{\partial x} \sim n_e T_e V_{th} / L_T. \]  

(4)

We report in this Letter that this is indeed realized by ambipolar constraint in a nearly collisionless plasma. In the case that the cooling spot is a perfect particle and energy sink (e.g., a black hole), which can be modeled by an absorbing boundary, \( q_{en} \) itself obtains the convective energy transport scaling, \( q_{en} \sim n_e V_{th} T_e \). With a radiative cooling mass, which can be modeled as a thermobath, the boundary of which recycles all particles across the boundary but clamps the temperature to a low value \( T_{in} \ll T_0 \), the cold electrons thus produced can restore the free-stream scaling for \( q_{en} \sim \alpha n_e v_{th} T_e \), but its spatial gradient is ever the cooling flow regime retains the convective energy transport scaling of Eq. (4). As the result, a robust cooling flow appears to aggregate mass towards the cooling spot.

Most interestingly, in such a cooling flow regime, the plasma thermal collapse comes in the form of propagating fronts that originate from the cooling spot with characteristic speeds. There are totally four (three) propagating fronts for the thermobath (absorbing) boundary: two of them propagate at speeds that scale with \( v_{th,e} \), so are named electron fronts, while the other two are ion fronts that propagate at speeds that scale with the local ion sound speed \( c_i \) (the last ion front disappears for the absorbing boundary). Fig. 1 illustrates the structure of the four fronts that propagate into a hot plasma for the thermal collapse with a thermobath boundary. It is important to note that cooling of a nearly collisionless plasma produces strong temperature anisotropy, so one must examine the collapse of \( T_i \) and \( T_\perp \) separately.

Cooling of \( T_e \) in a nearly collisionless plasma is primarily through free-streaming loss of suprathermal electrons satisfying \( v_e < -\sqrt{2 e (\Delta \Phi)_{max} / m_e} \) into the radiative cooling spot. Here \( (\Delta \Phi)_{max} \) is the maximum reflective potential in the plasma with \( \Delta \Phi = \Phi_e - \Phi_i(x) \) and the constant \( \Phi_i \) for the upstream plasma potential. The precooling zone bounded by the precooling front (PF) and the precooling trailing front (PTF) has \( T_e \) unchanged and \( V_e \approx 0 \), but a lowered \( T_i \), which is due to the depletion of fast electrons satisfying \( v_e \gg v_c = \sqrt{2 e (\Delta \Phi)_{max} - \Delta \Phi / m_e} \), yielding a truncated Maxwellian of the form

\[ f_e(v_{\parallel}, v_{\perp}) = \frac{n_m (\Phi(x))}{\sqrt{2\pi v_{th,e}^3}} e^{-v_{\parallel}^2 / 2v_{th,e}^2} \Theta \left( 1 - v_{\parallel} / v_c \right) \]  

(5)

with \( \Theta(1 - v_e / v_c) \) the Heaviside step function that vanishes for \( v_e > v_c \), and \( \delta(x) \) the Dirac delta function. The ambipolar electric field can draw some low-energy electrons to compensate for the loss of high-energy electrons and thus maintain quasi-neutrality. This in-falling cold electron population is modeled in Eq. (5) as a cold beam that due to ambipolar electric field acceleration has the speed \( v_e \). Between the PF and
PTF, the electron beam can be ignored, so
\[ T_{\parallel}(v_c) = \frac{\int m_e \tilde{v}_\parallel v f d\tilde{v}_\parallel}{\int f d\tilde{v}_\parallel} \approx T_0 \left[ 1 - \frac{v_c}{\sqrt{2\pi} v_{th,e}} e^{-\tilde{v}_\parallel^2/2v_{th,e}^2} \right], \tag{6} \]
where \( \tilde{v}_\parallel \equiv v_{\parallel} - v_{\parallel 0} \).

Eq. (6) predicts a detectable decrease in \( T_{\parallel}(v_c) \) from \( T_0 \) for \( v_c \approx 2.4v_{th,e} \), suggesting an electron PPF propagating at
\[ U_{PF} = 2.4v_{th,e}. \tag{7} \]
This corresponds to fast electrons with \( v_{\parallel} > U_{PF} \) traveling from the left boundary into the plasma, leaving behind a distribution at the PF with a void in \( v_{\parallel} > U_{PF} \). The PTF comes about due to the reflecting potential \((\Delta\Phi)_{RF} = (\Delta\Phi)_{max} - (\Phi_0 - \Phi_{RF})\) with \(\Phi_{RF}\) the ambipolar potential at the ion recession front (RF), which sets a lower cutoff speed \( v_c \) at
\[ U_{PTF} \equiv \sqrt{2e(\Delta\Phi)_{RF}/m_e}. \tag{8} \]
The deeper void now gives rise to a further reduced \( T_{\parallel} \),
\[ T_{\parallel}(U_{PTF}) \approx T_0 \left[ 1 - \sqrt{e(\Delta\Phi)_{RF}/\pi T_0 e^{-e(\Delta\Phi)_{RF}/T_0}} \right]. \tag{9} \]
The PTF rides these electrons that are reflected by the reflecting potential, and propagates at \( U_{PTF} < U_{PF} \). Since the ambipolar reflecting potential must satisfy \( e(\Delta\Phi)_{RF} \sim T_0 \) in a nearly collisionless plasma, \( U_{PTF} \approx v_{th,e} \) and \( T_{\parallel}(U_{PTF}) \) is only mildly cooler than \( T_0 \). Furthermore, \( T_{\parallel} \) and \( \Phi \) vary little between the RF and PTF, since the cutoff velocity remains the same at \( U_{PTF} \).

The ion flow remains vanishingly small ahead of the RF, so the electron cooling between the RF and PF is the result of electron conduction, which for the model \( f_e \) in Eq. (5) with \( v_c > v_{th,e} \) takes the form
\[ q_{en} \equiv \int m_e \tilde{v}_\parallel^3 f_e d\tilde{v}_\parallel \approx -\frac{n_m v_{th,e} T_0}{\sqrt{2\pi}} \left( \frac{v_c^2}{v_{th,e}^2} - 1 \right) e^{-v_c^2/2v_{th,e}^2} \tag{10} \]
\[ + n_b T_0 \frac{v_c^3}{v_{th,e}^3}. \]
Between PTF and PF, \( n_b \approx 0 \) and \( U_{PTF} \leq v_c \leq U_{PF} \) so \( q_{en} \) does scale as the free-streaming limit of Eq. (2), but with \( \alpha \) modulating in space as a function of \( v_c \). In fact, for \( v_c > \sqrt{2} v_{th,e} \), one finds
\[ \frac{dq_{en}}{dx} \approx n_m v_c \frac{\partial T_{\parallel}}{\partial x}, \tag{11} \]
so the solution of the energy equation, \( \partial T_{\parallel}/\partial t = -v_c \partial T_{\parallel}/\partial x \), reveals that \( v_c \) is the recession speed of \( T_{\parallel} \), reaffirming the particle picture noted earlier that the momentum space void in \( f_e \) propagates upstream with a speed of \( v_c \). This large \( q_{en} \) drives fast propagating electron fronts (PF and PTF) but produces modest amount of \( T_{\parallel} \) cooling for the large cutoff speed \( v_c = U_{PTF} \).

Much more aggressive cooling would need to occur as the plasma approaches the radiative cooling spot that is clamped at \( T_w \ll T_0 \). These are facilitated by the ion fronts that provide the reflecting potential \((\Delta\Phi)_{RF} \). The RF is where \( n_i \approx n_e \) starts to drop, and behind which plasma pressure gradient drives a cooling flow toward the radiative cooling spot. The main reflection potential, which is tied to the electron pressure gradient, is also behind the ion RF. An ion recession layer bounded by the RF and the cooling front (CF) is similar to the rarefaction wave formed in the cold plasma interaction with a solid surface [36 37], where the plasma parameters recede steadily with the local sound speed. What is different for the thermal quench of a nearly collisionless plasma is the large plasma temperature and pressure gradient and the nature of the heat flux. The electron flow associated with \( f_e \) in Eq. (5) within the recession layer is
\[ n_e V_{en}(v_c) = -\frac{n_m v_{th,e}^2}{\sqrt{2\pi}} e^{v_c^2/2v_{th,e}^2} + n_b v_c, \tag{12} \]
with \( n_e V_{en}(v_c) = \left( 1 + \text{Erf} \left( \frac{v_c}{\sqrt{2} v_{th,e}} \right) \right) n_m/2 + n_b \). For an absorbing boundary \((n_b = 0)\), a cutoff speed around \( v_{th,e}, v_c \approx v_{th,e} \sqrt{2\ln(v_{th,e}/V_{\parallel})} \), is sufficient to produce a \( V_{\parallel} \) that matches onto the increasing ion flow, \( V_{\parallel} \approx V_{\parallel0} \), for ambipolar transport through the recession layer. The in-falling cold electron beam reduces \( v_c \) and hence produces a lower reflecting potential across the recession layer as elucidated in Eq. (12).

The physics of \( q_{en} \) in the recession layer can be elucidated by rewriting Eq. (10) as
\[ q_{en} = \left( \frac{v_c^2}{v_{th,e}^2} + 2 \right) n_e V_{en} T_0 - 2n_b v_c T_0 - 3n_e V_{\parallel} \approx V_{\parallel} + 2n_m V_{\parallel0}^3. \tag{13} \]
For an absorbing wall, \( n_b = 0 \), and one finds \( q_{en} \) itself has a convective energy transport scaling: \( q_{en} \sim n_e V_{en} T_0 \). The condition for the cooling flow regime, Eq. (4), is obviously satisfied. In the case of a radiative cooling spot that produces copious amount of cold electrons, the leading order of \( q_{en} \sim -2n_b v_c T_0 \approx n_b v_{th,e} T_0 \) follows the free-streaming limit of Eq. (2). Remarkably the plasma thermal quench still produces a cooling flow, in which case Eq. (4) is satisfied due to the collisionless cold beam in the ambipolar electric field follows flux conservation \( n_b v_c = \text{constant} \), so \( \partial(-2n_b v_c T_0)/\partial x = 0 \) and the remaining terms in \( q_{en} \) have convective energy transport scaling. The VPIC [35] kinetic simulations shown in Fig. 2 confirms that convective scaling of Eq. (4) holds in the recession layer. In other words, electron cooling in a nearly collisionless plasma is modified by ambipolarity in such a way that large \( T_{\parallel} \) gradient can be supported in the recession layer to drive a cooling flow.

The propagation speed of the RF can be understood by ex-
amining the ion dynamics in the recession layer \[38,39\]

\[
\frac{\partial}{\partial t} n_i + \frac{\partial}{\partial x} (n_i V_{\parallel i}) = 0, \tag{14}
\]

\[
m_i n_i \left( \frac{\partial}{\partial t} V_{\parallel i} + V_{\parallel i} \frac{\partial}{\partial x} V_{\parallel i} \right) + \frac{\partial}{\partial x} (p_{\parallel i} + p_{\parallel e}) = 0, \tag{15}
\]

\[
n_i \left( \frac{\partial}{\partial t} T_{\parallel i} + V_{\parallel i} \frac{\partial}{\partial x} T_{\parallel i} \right) + 2n_i T_{\parallel i} \frac{\partial}{\partial x} V_{\parallel i} + \frac{\partial}{\partial x} q_{\parallel m} = 0, \tag{16}
\]

where we invoked the electron force balance \(e_n E_{\parallel} \approx -\partial p_{\parallel e}/\partial x\) and quasi-neutrality \(n_i = Z n_e\) with \(Z\) the ion charge, and \(p_{\parallel e} = n_e T_{\parallel e}\). Introducing a parameterization of \(q_{\parallel m} \approx \sigma_i n_i V_{\parallel i} T_{\parallel i}\), which is known from Ref. \[40\] and \(\partial q_{\parallel m}/\partial x \approx \sigma_i \partial (n_i V_{\parallel i} T_{\parallel i})/\partial x\) from Eq. (4), we obtain an universal length scale for \(p_{\parallel e}\)

\[
d\ln p_{\parallel e} \approx \mu \frac{d\ln \rho_{\parallel e}}{dx}, \tag{17}
\]

where \(\mu = (3 + \sigma_i)/(3 + \sigma_e) \times \left[-U + (1 + \sigma_i) V_{\parallel i}/[-U + (1 + \sigma_e) V_{\parallel i}]\right]. It is interesting to note that \(U > 0\) and \(V_{\parallel i} < 0\) have opposite sign in the recession layer where a cooling flow resides.

As a result, Eqs. (14-16) have self-similar solutions with similarity variable \(\xi = x - Ut\) with \(U\) being the local recession speed. We find

\[
U = \left[\sigma_i^2 V_{\parallel i}^2/4 + (1 + \sigma_i/3) c_s^2\right]^{1/2} + (1 + \sigma_e/2) V_{\parallel i}, \tag{18}
\]

where \(c_s = \sqrt{3(\mu Z T_{\parallel e} + T_{\parallel i})/m_i}\) is the local sound speed of a nearly collisionless plasma with anisotropic temperatures. At the ion recession front, \(V_{\parallel i} \approx 0\), so the speed of the ion recession front is

\[
U_{RF} = \sqrt{1 + \sigma_e/3c_s}. \tag{19}
\]

For \(Z = 1\), \(\sigma_i = 1\), \(\mu = 1\) and \(T_{\parallel e} \approx T_{\parallel i} = T_{0}\) at the recession front, we have \(U_{RF} \approx 2.8 v_{th,i}\) with \(v_{th,i} = \sqrt{T_{0}/m_i}\) the ion thermal speed, which agrees well with the simulation result. It is worth noting that the self-similar solution of Eq. (18) also recovers a known constraint \[41\] on the plasma exit flow at an absorbing boundary where a non-neutral sheath would form next to it as shown in the Supplement material.

In the absence of an absorbing boundary, the mass aggregated by the cooling flow will pile up, and the resulting back-pressure can now drive a second ion front (cooling front, CF). Behind the CF, \(T_{\parallel i}\) equilibrates with \(T_e\) as shown in Fig. 3. Such a deep cooling of \(T_{\parallel i}\) is through thermal conduction as indicated in Fig. 2. When the cooling flow runs into this nearly static plasma, the ion flow energy, which is substantial in the cooling flow, is converted into ion thermal energy via a plasma shock as shown in Fig. 3. Matching the conserved quantities across the shock while ignoring the heat flux, we find that the speed of the shock, which propagates upstream into the plasma, is simply the upstream sound speed at the shock front. The CF is the shock front, so its speed is

\[
U_{CF} = c_s(x = x_{CF}). \tag{20}
\]

Since the plasma temperature at the CF is considerably lower than that at the RF, we have \(U_{CF} < U_{RF}\). Generally, the colder \(T_e\), the smaller \(U_{CF}\). The presence of the CF and the cooling zone behind it, is of fundamental importance to \(T_{\perp e}\) and \(T_{\perp i}\) cooling as the cold particles provide dilutional cooling. It is also the source of cold electrons that are accelerated by the ambipolar electric field into the recession layer and beyond, cooling down \(T_{\perp e}\) further upstream.

In conclusion, the thermal collapse of a nearly collisionless plasma due to its interaction with a localized particle or energy sink, is associated with a cooling flow toward the cooling spot. This applies to unmagnetized plasmas, for example, in astrophysical systems, and magnetized plasmas, for example, in earth’s magnetosphere or a tokamak fusion plasma. It is the fundamental constraint of ambipolar transport, along the field line in a magnetized plasma, that limits the spatial gradient of electron (parallel) heat flux to the much weaker convective \((V_{\parallel e})\) scaling as opposed to the free-streaming \((v_{th,e})\) scaling. Such weaker scaling is essential to sustain a temperature and hence pressure gradient for driving the cooling flow toward
the cooling spot over the ion recession layer. The cooling flow eventually terminates against the cooling spot via a plasma shock that converts the ion flow energy into ion thermal energy. This shock or cooling front propagates away from the cooling spot at upstream ion sound speed, and it has the most profound role in the deep cooling of the surrounding hot plasmas, especially the ions. Unlike the ions, the electrons can be cooled ahead of the recession front due an electron heat flux that follows the free-streaming limit (\( q_{\text{en}} \propto n_e v_{\text{thi}} T_e \)). Interestingly this large heat flux does not imply significant cooling of \( T_{e\parallel} \) in a nearly collisionless plasma ahead of the recession front, but induces a very limited amount of \( T_{e\parallel} \) drop over a very large volume, because the precooling and precooling trailing fronts have propagation speeds that scale with electron thermal speed.

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