Staged cooling of a fusion-grade plasma in a tokamak thermal quench

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Abstract

In tokamak disruptions where the magnetic connection length becomes comparable to or even shorter than the plasma mean-free-path, parallel transport can dominate the energy loss and the thermal quench of the core plasma goes through four phases (stages) that have distinct temperature ranges and durations. The main temperature drop occurs while the core plasma remains nearly collisionless, with the parallel electron temperature $T_e\parallel$ dropping in time as $T_e\parallel \propto t^{-2}$ and a cooling time that scales with the ion sound wave transit time over the length of the open magnetic field line. These surprising physics scalings are the result of effective suppression of parallel electron thermal conduction in an otherwise bounded, quasineutral, and collisionless plasma, which is different from what are known to date on electron thermal conduction along the magnetic field in a nearly collisionless and quasineutral plasma.

Keywords: staged cooling, thermal quench, convection, conduction, tokamak disruption

The energy confinement time of a fusion-grade plasma in an ITER-like tokamak reactor is around 1 second (s) \cite{1}. This is set by plasma transport across magnetic field lines, or more accurately, nested magnetic surfaces \cite{2}. When a major disruption occurs, such a plasma is expected to lose most of its thermal energy over a period of around 1 millisecond (ms) or so, as noted in the original ITER design document \cite{3}. A dedicated experimental campaign on JET \cite{4}, after the initial ITER design was completed \cite{5}, revealed that the thermal quench (TQ) duration could vary greatly, from 0.05 to 3.0 ms. The significance of a tokamak TQ is that it marks the point of no return in a tokamak disruption. It not only brings a thermal load management issue at the divertor plates and first wall, but also determines the runaway seeding for the subsequent current quench as well as the parallel electric field that could drive avalanche growth of runaway electrons. Since the faster a TQ is, the more problematic it becomes for effective mitigation, there are practical interests and urgencies in understanding the fundamentals of transport physics that are responsible for the rapid TQ, particularly the large variation in its duration.

The 3–4 order of magnitude faster plasma energy loss rate in a tokamak disruption \cite{4, 6, 7} is intuitively understood as the result of parallel transport along open magnetic field lines dominating over the perpendicular transport. For that to materialize, the magnetic field lines must connect the hot fusion plasma of 10–15 KeV temperature directly to a cold and dense plasma, which would also serve as a radiative energy sink. This is known to occur in at least two scenarios. The first and more common one of the naturally occurring disruptions has the globally stochastic magnetic field lines connecting the hot core plasma directly onto the divertor surface and/or the first wall,
as the result of large-scale magnetohydrodynamic instabilities destroying nested magnetic surfaces [8–11]. The other is high-Z impurity injection in the form of deliberately injected solid pellets or accidentally falling tungsten debris [12–15]. In both situations, a nearly collisionless plasma is made to intercept a radiative cooling mass, being that an ablated pellet or a vapor-shielded wall.

The simplest problem setup to decipher the parallel cooling physics is to unwind the open field lines into a slab and have a hot plasma of temperature $T_0$ bounded by thermbath boundary at the two ends that recycles plasma particles but clamps the temperature of the recycled plasma particles to $T_w \ll T_0$. This setup is studied in this work by first-principles 1D3V VPIC [16] simulations that fully account for Coulomb collisions. The length of the slab or field line, $x \in [-L_B, L_B]$, is twice the magnetic connection length $L_B$. The other characteristic length of the problem is the parallel mean-free-path $\lambda_{\text{mfp}}$, which is tens of kilometers for a plasma of $T_0 = 10–15 \text{ KeV}$ and $n_e = 10^{19–20} \text{ m}^{-3}$. The ratio of the two is the Knudsen number $K_n \equiv \lambda_{\text{mfp}}/L_B$, and it sets the different cooling regimes (stages) that a TQ can go through. For an initial plasma of $K_n \geq 1$, the TQ would start with the (nearly) collisionless phase (regime) and eventually transition to the collisional phase after the plasma cools down sufficiently as $K_n \sim T_e^2/n_e$ [17–20]. The simplicity of this prototypical problem setup reinforces the general importance and broad applicability of the underlying plasma cooling physics in fusion as well as non-fusion applications, which will be shown in this paper to present a number of surprises.

The prevailing view on plasma TQ is the dominant role of electron parallel thermal conduction. This is described in the collisional limit by Braginskii closure for parallel electron conduction flux [21]

$$q_{\|} = \frac{1}{2} \int m_e \tilde{v}_\| [f_e d^3 v = -3.16 \frac{n_e T_e \tau_e}{m_e} \frac{\partial T_e}{\partial x} \sim n_e \tilde{v}_{th,e} T_e K_n, \quad (1)$$

with $f_e$ the full electron distribution function, $\tau_e$ the electron collision time, $m_e$ the electron mass, $\tilde{v}_{th,e} = \sqrt{T_e/m_e}$ the electron thermal speed. As the collisionality reduces, $q_{\|}$ scales up linearly with $K_n$. $q_{\|} \propto K_n$. This would start to break down when $K_n$ gets above $10^{-2}$ or so, and when $K_n \sim 1$ or $K_n \gg 1$, $q_{\|}$ would retain the $n_e \tilde{v}_{th,e} T_e$ scaling but the $K_n$ term is replaced by a saturated numerical factor $\alpha \sim 0.1$ [18],

$$q_{\|} = c n_e \tilde{v}_{th,e} T_e, \quad (2)$$

With the so-called flux-limiting form of equation (2) in the collisionless or free-streaming regime [22], one finds the solution of the heat conduction equation in the bounded domain of $x \in [-L_B, L_B]$ showing a TQ duration

$$\tau_{\text{TQ}} \sim L_B/\tilde{v}_{th,e}. \quad (3)$$

This suggests a very fast TQ indeed, on the order of the thermal electron transition time over the open magnetic field line, for a tokamak plasma having $K_n \sim 1$ or $K_n \gg 1$ at the onset of a disruption.

In this paper, we show that the core plasma cooling in a bounded plasma like that of a tokamak disruption follows qualitatively and quantitatively different behaviors. Specifically the parallel electron temperature $T_{\|} = \int m_e \tilde{v}_\| [f_e d^3 v \sim f_e d^3 v$ goes through four distinct stages with durations $\Delta t_{1,2,3,4}$ segmented by transition temperatures $T_{1,2,3,4}$, as illustrated in figure 1. There is a precursor stage of very short duration at the electron transit time $\Delta t_1 \sim \tau_e^2 \equiv L_B/\tilde{v}_{th,e}$, the initial electron thermal speed, which follows the same scaling as in equation (3). Interestingly it produces rather limited cooling of the core plasma, $T_1 \geq 0.6T_0 \gg T_w$. The core plasma cooling is primarily accomplished in the next collisionless stage from $T_1 \sim T_0$ to $T_2 \sim T_0$, with $\Delta t_2 \gg \Delta t_1$. There is a transition stage $\Delta t_2 \sim \Delta t_2$ that connects the collisionless cooling phase to the eventual collisional cooling phase, in which $\Delta t_3$ is far longer than the main collisionless cooling stage of $\Delta t_2$. The most surprising and impactful finding is that the main collisionless cooling stage from $T_1 \sim T_0$ to $T_2 \ll T_0$, is dominated by convective energy transport as opposed to the commonly accepted much faster electron parallel thermal conduction, e.g. equation (2). This yields a TQ time $\Delta t_2$ that scales with the ion transit time $\Delta t_2 \propto \tau_{tr}^2 \equiv L_B/c_i$ with $c_i = \sqrt{2/(1+Z)T_0/m_i}$ the ion sound speed ($Z$ is the ion charge). The main TQ time of $\Delta t_2$ is thus qualitatively different from that of $\tau_{\text{TQ}}$ in equation (3), which constitutes the prevailing textbook description of thermal loss by electron free-streaming along magnetic field lines.

**Precursor phase ($T_{\|}$ from $T_0$ to $T_1$ and $\Delta t_1 \sim \tau_e^2$):** At the onset of a plasma TQ when the hot plasma on the open...
field lines suddenly intercepts a cold boundary, $T_{\parallel}(x = 0, t)$ barely changes over a period that is approximately one-third of $\tau_{en}$. This is simply the time for the so-called precooling front (PF), which is driven by suprathermal electron loss and propagates from the boundary toward the core plasma at a speed $U_{PF} = 2.4 v_{th}$, to arrive at the plasma center $(x = 0)$ [23]. This instance is marked by the black dashed vertical line in figure 2. The subsequent cooling of $T_{\parallel}(x = 0)$ is indeed driven by the electron conduction flux that follows a similar scaling as the flux limiting form in equation (2). Notice that for a nearly collisionless plasma, the conduction flux corresponding to $T_{\parallel}$ cooling is given by $q_{en} \equiv \int m_{e} \nu_{e} d^{3}v$. Consequently, the duration of this phase does follow the free-streaming scaling previously given in equation (3). Remarkably the precursor phase ends abruptly when the so-called precooling trailing front (PTF) reaches the center. This is a second electron front that propagates from the boundary toward the core plasma, with a speed $U_{PF} = \sqrt{2 e \Delta \Phi_{RF}/m_{e}}$ [23]. Here the reflecting potential $\Delta \Phi_{RF}$ is the result of the ambipolar electric field in the ion recession layer where a plasma rarefaction wave is formed. Since the ambipolar potential scales with $T_{\parallel}$, we would normally have $v_{th, e} < U_{PF} < U_{PF}$, and thus $\Delta t_{\parallel} \sim \tau_{en}$. The short duration of the precursor phase produces very limited cooling and typically $T_{\parallel} \geq 0.6 T_{0}$.

Cooling flow phase ($T_{\parallel}$ from $T_{1}$ to $T_{2}$ with $T_{w} < T_{2} \ll T_{0}$ and $\Delta t_{\parallel} \sim \tau_{en}$): Once the PTFs reach the plasma center from both ends, $q_{en}$ undergoes a qualitative transition from scaling with $v_{th, e}$ (i.e. $q_{en} \propto n_{e} v_{th, e} T_{\parallel}$) to scaling with $V_{\parallel}$ (i.e. $q_{en} \propto n_{e} V_{\parallel} T_{\parallel}$). Here $V_{\parallel}$ is the parallel ion convective flux and ambipolar transport implies the parallel electron flow $V_{\parallel} \approx V_{\parallel, \text{th}}$. Previous we have shown that in a semi-infinite plasma bounded at one side only by a thermobath boundary, $q_{en}$ always has the free-streaming flux-limiting form $n_{e} v_{th, e} T_{\parallel}$ due to the cold electrons being pulled into the hot plasma by the ambipolar electric field [23]. The qualitative change of $q_{en}$ in a bounded plasma is driven by how $T_{\parallel}$ is being cooled in a long mean-free-path plasma where electron trapping is provided by the reflecting (ambipolar) potential along the magnetic field line on both ends.

The core cooling physics is most straightforwardly understood in a plasma with perfectly absorbing boundaries. Since the trapped-passing boundary is given by $\nu_{c} \equiv \sqrt{e \Delta \Phi_{RF}/m_{e}}$, and the truncated electron distribution at $x = 0$ has the form

$$f_{0} = \frac{n_{0}}{(2\pi)^{3/2} v_{th, 0}^{3}} \left\{ e^{-\nu_{i}^{2} + v_{i}^{2}} \right\} / 2 v_{th, 0} \Theta \left( 1 - \frac{V_{\parallel}}{v_{c}} \right) \Theta \left( 1 + \frac{V_{\parallel}}{v_{c}} \right)$$

with $\Theta(x)$ the Heaviside function satisfying $\Theta(x < 0) = 0$ and $\Theta(x > 0) = 1$. One can see that the parallel electron temperature $T_{\parallel}(\nu_{c}, T_{0}) = \left\{ m_{e} \nu_{e}^{2} v_{th, 0}^{2} / \int f_{0}^{(1)} d^{3}v \right\} = (\nu_{i}^{2} / 2 v_{th, 0}^{2}) T_{0}$ of the electrostatically trapped long mean-free-path electrons cools drastically with a decreasing $v_{c}/v_{th, 0} \ll 1$. This correlates with a reducing reflecting potential $\Delta \Phi_{RF}$. Away from the symmetry point at $x = 0$, a finite parallel electron conduction flux arises from the asymmetric truncations ($V_{L}$ and $V_{R}$) on the positive and negative sides of electron distribution in $V_{\parallel}$

$$f_{\pm} = \frac{n_{0}}{(2\pi)^{3/2} v_{th, 0}^{3}} \left\{ e^{-\nu_{i}^{2} + v_{i}^{2}} \right\} / 2 v_{th, 0} \Theta \left( 1 - \frac{V_{\parallel}}{v_{c}} \right) \Theta \left( 1 + \frac{V_{\parallel}}{v_{c}} \right)$$

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In the limit of significant plasma cooling, one can write $V_{\parallel} \approx (-V_{L} + V_{R}) / 2$ and find a convective energy flux scaling for electron thermal conduction

$$q_{en} \approx \frac{6 T_{\parallel}}{5 T_{0}} n_{e} V_{\parallel} T_{\parallel}$$

In a cooling plasma of $T_{\parallel} \ll T_{0}$, one recovers the remarkable result of $q_{en} \ll n_{e} V_{\parallel} T_{\parallel}$ observed in VPIC simulations.

An even more interesting and impactful result is on $q_{en}$ of a plasma with thermobath boundary conditions. The subtlety is that cold electrons from the boundary plasmas can follow the ambipolar electric field into the core plasma. This cold electron beam component was previously found to restore the ambipolar electric field into the core plasma. This cold electron beam component was previously found to restore the ambipolar electric field into the core plasma.
the cooling flow $V_{\parallel}$ linearly grows in $x$ from zero at $x = 0$ to the Bohm speed $u_{\text{Bohm}}$ near $x = L_B$, and it decays in time due to plasma cooling. This can be understood via a separable solution $V_{\parallel}(x,t) = P(x)Q(t)$ to the ion momentum equation, where the decay of the plasma flow is mostly balanced by the inertial term (e.g. see figure 3),

$$\frac{\partial}{\partial t} V_{\parallel} + V_{\parallel} \frac{\partial}{\partial x} V_{\parallel} = -\nabla p \approx 0.$$  

(5)

The separable solution has

$$dP/dx = -dQ/dt = C_0,$$  

(6)

with $C_0$ a constant. The solution $P(x) = C_0 x$ underlies the aforementioned VPIC simulation result. In the normal situation with an upstream Maxwellian plasma, $u_{\text{Bohm}} \approx \sqrt{(Z/\ell_{\parallel}) + 3T_{\parallel}}/m_i$ because of the parallel heat flux into the sheath [24, 25]. While for the nearly collisionless plasma in the cooling flow phase during a TQ, the thermal conduction is insignificant, so $u_{\text{Bohm}} \approx \sqrt{3ZT_{\parallel} / m_i}$. Since the cooling starts at $T_0$, the Bohm speed at $t = 0$ is simply the parallel sound speed $c_s \equiv \sqrt{3(Z+1)T_0 / m_i}$ in a magnetized plasma with strong temperature anisotropy, so we can write $C_0 = c_s / L_B$, which introduces an ion transit time $\tau_{tr}^i = L_B / c_s$ for $Q$,

$$Q = \left(\delta + \frac{t}{\tau_{tr}^i}\right)^{-1}.$$  

(7)

Here $\delta \approx 1$ depends on the condition when $T_{\parallel}$ becomes uniform and we assumed $T_{\parallel} \approx T_{\|}$ due to the convective heat flux dominating the cooling. Similarly a separable solution $T_{\parallel} = M(x)N(t)$ can be found from the electron energy equation [26],

$$n_e \frac{\partial}{\partial t} T_{\parallel} + n_e V_{\parallel} \frac{\partial}{\partial x} T_{\parallel} + 2n_e T_{\parallel} \frac{\partial}{\partial x} V_{\parallel} + \frac{\partial}{\partial x} q_{\perp} = 0.$$  

(8)

Ignoring the subdominant $\partial q_{\perp} / \partial x$ term and the $n_e V_{\parallel} \partial T_{\parallel} / \partial x$ term because of weak $T_{\parallel}$ variation, one finds $d\ln N(t)/dt = -2C_0 Q(t)$ for

$$\frac{T_{\parallel}(t)}{T_0} = N(t) = \left(\delta + \frac{t}{\tau_{tr}^e}\right)^{-2},$$  

(9)

with $M(x) \approx T_0$. This reveals two pieces of interesting and important physics: (1) the characterized time of collisionless core $T_{\parallel}$ cooling is given by the ion sound transit time $\Delta t_i \sim \tau_{tr}^i$, and (2) $T_{\parallel}(t) \propto t^{-2}$ for $t/\tau_{tr}^e \gg 1$ as shown in figure 2. Substitution of equation (9) into $u_{\text{Bohm}}$ gives $u_{\text{Bohm}} = P(x = L_B)Q$ as expected so that the solutions are self-consistent.

The temperature $T_2$ signifies the beginning of a transition from collisionless to collisional cooling, and it is set by the constraint that the convection flux equals the Braginskii conduction, which suggests a Knudsen number at the transition

$$K_{n,2} \sim c_s/v_{th,e} \sim \sqrt{m_e(1+Z)/m_i}.$$  

(10)

If we ignore the small density variation at the core (see figure 3) so that $K_{n,2} \approx K_{n,0}(T_2 / T_0)^2$, we find

$$T_2 \approx T_0 K_{n,0}^{-1/2} \left[m_e(1+Z)/m_i\right]^{1/4}.$$  

(11)

It follows that $T_2 \approx 0.04T_0$, agreeing well with VPIC simulation data in figure 2. The cooling time from $T_1$ to $T_2$ is then found from equation (9) to be

$$\Delta t_2 \sim K_{n,0}^{1/4} \left[m_e(1+Z)/m_i\right]^{-1/8} \tau_{tr}^i.$$  

(12)

For plasma parameters in figure 2, $\Delta t_2 \sim 5\tau_{tr}^i$, in good agreement with the simulation data. 

*Transition phase* $(T_{\parallel} \sim T_2 \sim T_3$ and $\Delta t_3 \sim \tau_{tr}^i)$: The cooling flow meets its eventual collapse when the boundary plasma becomes over-pressured because of density build-up in the presence of decreasing temperature gradient. This sets off a recompression of the core that transiently heats the plasma temperature, especially that of the ions. Due to the strong Landau damping of ion-acoustic waves, the recompression flow is removed over one thermal ion transit period, so $\Delta t_3 \approx \tau_{tr}^i$. This $\Delta t_2$, leaving a core plasma temperature $T_3 \approx T_2$ at the end of the one-bounce transition period, as shown in figure 2. It should be noted that (1) the recompression
\(T_{\epsilon}\) bump at \(\tau_3\) can be minimal depends on initial plasma condition and the magnetic connection length; and (2) such a bulk temperature bump has not yet been observed experimentally.

Collisional cooling phase \((T_{\epsilon})\) from \(T_3\) to \(T_e\), and \(\Delta t_4 \gg \tau_{\epsilon}^4\): Collisional cooling is dominated by the thermal conduction \(q_{\|}\) given by Braginskii in equation (1):

\[
3n_e \frac{\partial}{\partial t} T_{\epsilon\|} + \frac{\partial}{\partial x} q_{\|} = 0. \tag{13}
\]

A separable solution \(T_{\epsilon\|} = M(x)N(t)\) with \(M(x) \neq \text{const.}\) has the core \(T_{\epsilon\|}\) evolve in time as

\[
\frac{T_{\epsilon\|}}{T_3} = \left\{ \delta_2 + 1.6 K_{n,3} \sqrt{\frac{T_{\epsilon\|}}{m_e L_B}} \right\}^{-2/5}, \tag{14}
\]

where \(\delta_2 \sim 1\) accounts for \(\Delta t_2\) and \(\Delta t_3\) in \(t\). Here all the variables are normalized using the quantities at the beginning of collisional TQ including the Knudsen number \(K_{n,3}\). Therefore, this solution can be applied to plasma that is initially within the collisional regime. For an initially collisionless fusion plasma, we can substitute \(T_3 \approx T_2\) and hence \(K_{n,3} \approx K_{n,2}\) into equation (14) to obtain

\[
\frac{T_{\epsilon\|}}{T_3} = \left\{ \delta_2 + 1.6 K_{n,0}^{1/4} \left[ \frac{m_e (1 + Z)}{m_i} \right]^{1/8} \frac{t}{\tau_{\epsilon}^4} \right\}^{-2/5}. \tag{15}
\]

Here \(t\) is naturally normalized by \(\tau_{\epsilon}^4\) as in equation (9) for the cooling flow phase. The factor in front has weak powers \((-1/4\) and \(1/8\) respectively) of \(K_{n,0}\) and \(m_e (1 + Z)/m_i\), so it is also a number of order unity, like that in equation (9). What is different is that \(T_{\epsilon\|} (t) \propto t^{-2/5}\) as shown in figure 2, in sharp contrast to \(T_{\epsilon\|} (t) \propto t^{-2}\) in the cooling flow phase. This explains a much slower core TQ in the collisional regime than that in the collisionless regime. Specifically, the core TQ from \(T_3\) to \(T_e\) takes

\[
\Delta t_4 \sim \left( \frac{T_0}{T_e} \right)^{-5/2} K_{n,1}^{-1} \left[ \frac{m_e (1 + Z)}{m_i} \right]^{1/2} \tau_{\epsilon}^4. \tag{16}
\]

It is easy to check that \(\Delta t_4 \gg \Delta t_5\) for \(T_e < T_2\).

In conclusion, the main core thermal collapse \((T_3 \rightarrow T_2 < T_0)\) is mostly accomplished in the cooling flow phase with a duration that is a few times the ion sound transit time \(\Delta t_2 \sim L_B/c_s\), due to the surprising physics that electron thermal conduction is subdominant to convective energy transport in a nearly collisionless cooling plasma. This is followed by a transition phase in which over-pressured boundary plasma reheats the core plasma by compression. The final collisional cooling phase takes so long that the experimentally observed deep cooling to tens or a few eVs in milliseconds or shorter time must be due to a different mechanism, with impurity radiation a leading candidate, notwithstanding the potential inconsistency with the time scale for impurity transport from the boundary to the core. The physics scaling reveals that for a fusion-grade plasma in which \(T_0\) and \(n_0\) are severely constrained, the variability in TQ duration, which experimentally is associated with \(\Delta t_{3,4}\), is mostly set by the magnetic connection length \(L_B\). This potentially provides a means to determine the otherwise difficult-to-access \(L_B\) of globally stochastic magnetic fields. To quantify our findings, we show \(\Delta t_{3,4}\) from equations (12) and (16) of representative DIII-D and ITER plasmas for a range of \(L_B\) in figure 4. This provides a baseline prediction of TQ duration for comparison with experimental measurements and more involved transport calculations.

A key unresolved issue from the 1D3V calculations shown here is the role of perpendicular transport, both convective and diffusive. Of particular interest are (1) the \(E \times B\) mix driven by the ambipolar electric field in 3D magnetic fields, which was previously shown in collisionless 3D gyrokinetic simulations [27, 28]; (2) the Rechester–Rosenbluth effect due to stochastic magnetic field lines, which is known to be particularly effective in weak stochastic magnetic fields [29]; and (3) turbulent fluctuations induced cross field transport, which is observed to have a significant role if the field line stochasticity is weak in \(B_{\perp}+\) simulations [30]. It is of great interest and importance to perform self-consistent 3D simulations that can establish how the perpendicular transport in 3D would impact the parallel transport physics observed in 1D3V calculations, especially in the weak magnetic chaos regime.

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