The First 3D Coronal Loop Model Heated by MHD Waves against Radiative Losses

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Abstract

In the quest to solve the long-standing coronal heating problem, it was suggested half a century ago that coronal loops could be heated by waves. Despite the accumulating observational evidence of the possible importance of coronal waves, still no 3D MHD simulations exist that show significant heating by MHD waves. Here we report on the first 3D coronal loop model that heats the plasma against radiative cooling. The coronal loop is driven at the footpoint by transverse oscillations, and subsequently the induced Kelvin–Helmholtz instability deforms the loop cross section and generates small-scale structures. Wave energy is transferred to smaller scales where it is dissipated, overcoming the internal energy losses by radiation. These results open up a new avenue to address the coronal heating problem.

Unified Astronomy Thesaurus concepts: Magnetohydrodynamics (1964); Magnetohydrodynamical simulations (1966); Solar coronal heating (1989); Solar coronal waves (1995)

Supporting material: animations

1. Introduction

The mechanisms to heat the corona to its temperature of several million kelvin have been divided between direct current (DC) and alternating current (AC) ones. The DC heating mechanisms are models that generate currents and magnetic reconnection. Given that previous detections of coronal waves were sporadic events (Aschwanden et al. 1999; Nakariakov et al. 1999), modeling efforts have mainly gone into 3D DC heating models (Gudiksen & Nordlund 2005; Chen et al. 2014; Carlsson et al. 2016). However, the recent observational evidence for the ubiquitous presence of transverse coronal waves (Tomczyk et al. 2007; McIntosh et al. 2011; Anfinogentov et al. 2015) motivates a revisiting of the AC heating mechanisms. Wave heating has been successfully achieved in 1D (Suzuki & Inutsuka 2005) and reduced 3D MHD (van Ballegooijen et al. 2011, 2014) models. Recently, Matsumoto (2018) extended the 1D wave heating models to 3D, but the resulting model lacked any transverse structure and thus the waves heat a flux tube without a density-enhanced coronal loop. Recent 3D models with density-enhanced coronal loops have not been able to sustain the corona (Pagano et al. 2018; Pagano & De Moortel 2019). This is because the wave heating is insufficient to compete against the strong radiative cooling of the coronal loops. Since the optically thin radiative cooling directly decreases the internal energy, it is the key quantity to overcome in coronal heating models.

The observed transverse waves that show little to no decay (Tian et al. 2012; Wang et al. 2012) are called decayless kink waves (Nistico et al. 2013). Even though it is still unconfirmed whether these transverse waves have enough energy to heat the corona (Tomczyk et al. 2007; McIntosh et al. 2011; Van Doorselaere et al. 2014), the observed energy flux of transverse oscillations is usually underestimated because some “dark” energy is hidden in the nonthermal broadening of spectral lines (De Moortel & Pascoe 2012; McIntosh & De Pontieu 2012; Pant et al. 2019). The decayless waves are interpreted as being driven at the footpoints by photospheric motions, but the wave energy is transferred and dissipated in the corona, through either resonant absorption (Hollweg & Yang 1988; Goossens & Poedts 1992), mode coupling (Pascoe et al. 2012), phase mixing (Pagano & De Moortel 2017), or Kelvin–Helmholtz instability (KHI, Browning & Priest 1984; Terradas et al. 2008; Antolin et al. 2015). The last mechanism cascades the driving energy from large scales to small scales, where it is dissipated at low Reynolds numbers (Karampelas et al. 2017, 2019a; Afanasyev et al. 2019; Guo et al. 2019a). The resulting loops have a fully turbulent cross section (Karampelas & Van Doorselaere 2018; Antolin & Van Doorselaere 2019) and show decayless oscillations resembling observed events (Karampelas et al. 2019b).

In this paper, we study the heating effects of transverse oscillations on coronal loops via three-dimensional MHD simulations. Our work is different from previous studies in the following three aspects. (1) We explicitly include the radiative cooling in simulations. (2) Our results show that the dissipation of transverse oscillations is promising to balance the radiative cooling of the coronal loops. (3) As far as we know, we propose a new method estimating numerical viscosity and resistivity in MHD simulations. The numerical model is presented in Section 2. Section 3 provides the numerical results, and Section 4 is the discussion and conclusion.

2. Numerical Model

We model the coronal loop as a straight, density-enhanced magnetic flux tube in Cartesian coordinates. Gravity is ignored. The loop is along the z direction, with all physical parameters depending only on the transverse position x and y. Specifically, density is given by the relation

\[ \rho = \rho_1 + (\rho_2 - \rho_1) \zeta(x, y), \]

\[ \zeta(x, y) = \frac{1}{2}(1 - \tanh(b(\sqrt{x^2 + y^2}/R - 1))), \]  

(1)

where \([\rho_1, \rho_2] = [3 \times 10^8, 10^8] \text{cm}^{-3}\) denote the density in the interior and exterior of the loop, and \(R = 1 \text{Mm}\) is the loop
radius. $b$ is a parameter controlling the width of the transition layer $l$, which we choose as $b = 20$, corresponding to $l \approx 0.3 R$. To study the heating of the loop we set temperatures inside and outside the loop to be equal, $T_i = T_o = 1$ MK, in order to avoid a mixture of plasma at different temperatures. The magnetic field is along the $z$ direction, with the magnetic field strength inside the loop being $B_z = 30$ G. Stability is ensured by maintaining the total pressure balance. The Alfvén speed inside the loop is $v_{Ai} \approx 3778$ km s$^{-1}$.

Similar to previous works (Pascoe et al. 2010), we trigger a kink wave by implementing a time-dependent transverse velocity driver at the footpoint. The velocity inside the loop is

$$[v_x, v_y] = \left[ v_0 \cos\left(\frac{2\pi}{P_k}\right), 0 \right],$$

and the velocity outside the loop is

$$[v_x, v_y] = v_0 \cos\left(\frac{2\pi}{P_k}\right) R^2 \left[ \frac{x^2 - y^2}{(x^2 + y^2)^2}, \frac{2xy}{(x^2 + y^2)^2} \right],$$

where $P_k$ is the period of the fundamental kink mode. We choose the loop length $L = 200$ Mm, so that $P_k = 2L/c_k = 86.4$ s, with $c_k \approx 4672$ km s$^{-1}$ being the kink speed. In our simulations we set $v_0 = 8$ km s$^{-1}$, corresponding to the lower limits of chromospheric observations of the transverse waves (De Pontieu et al. 2007), which is larger than previously used typical values at the photosphere (e.g., Karampelas et al. 2017).

The simulations are performed using the ideal MHD module of the PLUTO code (Mignone et al. 2007). PLUTO solves the conservative form of the ideal MHD equations (i.e., continuity, momentum, induction, and energy equations). We use the finite-volume piecewise parabolic method for spatial reconstruction, and the second-order Runge–Kutta for time-stepping. The HLLD approximate Riemann solver is used for flux computation, and a hyperbolic divergence cleaning method is used to keep the magnetic field divergence-free. The simulation domain is $[-8, 8]$ Mm, $[-8, 8]$ Mm, and $[0, 200]$ Mm in $x$, $y$, and $z$ directions. To resolve the small-scale structures as far as possible in the transverse direction, we use 800 uniform grid points in both the $x$ and $y$ directions, and 100 in the $z$ direction, corresponding to a spatial resolution in the $x$ or $y$ direction of 20 km. For the side boundaries ($x$ and $y$), we use outflow boundary conditions for all quantities. The transverse velocities at the bottom boundary are given by the driver, and those at the top boundary are fixed to be zero. At both the top and bottom boundaries, the $z$-component of velocity is antisymmetric, while other quantities have zero-gradient boundary conditions. In the simulations, we have ignored resistivity, viscosity, and thermal conduction. The energy dissipation in our simulation is entirely caused by numerical dissipation at the grid scales.

### 3. Results

To study the heating effects of footpoint-driven kink waves in the presence of radiative cooling, we perform an MHD simulation with the radiative cooling included (radiative simulation hereafter). As the comparison case, an ideal MHD simulation (adiabatic simulation hereafter) similar to previous works (Karampelas et al. 2017; Guo et al. 2019a) is also conducted. The setup of the radiative simulation is the same as that of the adiabatic simulation initially, but with radiative cooling turned on when $t > 600$ s. The radiative cooling function $\Lambda(T)$ is tabulated from the Chianti database (Dere et al. 2019). At the same time as radiative cooling is turned on, in order to maintain the background corona, we also turn on a temporally and spatially independent constant heating term across the whole domain $H_0 = n_e^2 \Lambda_0$, where $n_e = 10^6$ cm$^{-3}$ and $\Lambda_0 = 2.6 \times 10^{-25}$ erg cm$^{-3}$ s$^{-1}$. $H_0$ is implemented to maintain the background corona by balancing its radiative loss. By having both the radiative cooling and the background heating, we are in fact analyzing the heating effects after the KHI has fully developed and only inside the density-enhanced coronal loops.

#### 3.1. Loop Dynamics

For both cases, we run simulations until $t = 2000$ s ($t \approx 23 P_k$). Figure 1 displays the 3D density structure of the subvolume with $\rho \geq (\rho_c + \rho_i)/2$, the temperature structure of the subvolume with $T \geq 1.1$ MK, and the forward-modeled intensity of the Fe XII 193.5 Å line by FoMo code (Van Doorsselaere et al. 2016) for the radiative simulation at $t = 860$ s ($t \approx 10 P_k$). From Figure 1 and the related movies we see the formation of a standing kink mode as the result of the footpoint driver. The loop is distorted first at the boundary due to the velocity shear, and then in its whole cross section as KHI develops. Heating is initiated at the loop boundary, and is more profound at the loop boundary than in the loop interior, due to strong velocity shear at the boundary and strong radiative loss in the loop interior. Heating occurs first at the footpoint and only later at the loop apex, creating the intensity enhancement of the Fe XII 193.5 Å line (formation temperature $\sim 1.58$ MK) first at the footpoint and then at the loop apex.

Figure 2 shows time–distance maps of density at the mid-plane (a slice at $y = 0$) of the loop apex for two simulation cases. We find that the introduction of radiative cooling does not have any obvious influence on the evolution of the wave amplitude. The average displacement of the loop center is around 0.6–0.7 Mm in our simulations, which is comparable to the larger amplitude limit of the observed decayless oscillations ($\sim 0.5$ Mm from Nisticò et al. 2013 and Anfinogentov et al. 2015).

In order to examine the loop dynamics in more detail, we plot in Figure 3 the density $\rho$, temperature $T$, $z$-component of current $J_z$, and $z$-component of vorticity $\omega_z$ at the loop apex ($z = 0.5 L$) and near the footpoint ($z = 0.1 L$) for both adiabatic and radiative simulations at $860$ s ($t \approx 10 P_k$). From Figure 3 and the related movie, we find different evolutions at the apex and near the footpoint. At the apex, the KHI is initiated at the loop boundary due to the strong velocity shear, then it develops further to the extent where the loop cross section is fully deformed. Near the footpoint, where mixing is less strong, the loop cross section is less deformed than at the apex. As simulations go on, both $J_z$ and $\omega_z$ develop. The high-temperature regions coincide with the regions of strong current or viscosity. $J_z$ is much stronger near the footpoint, while $\omega_z$ is more prominent around the apex, indicating that numerical resistive heating is stronger near the footpoint, while numerical viscous heating is more prominent around the apex (Van Doorsselaere et al. 2007; Karampelas et al. 2017). Some high-temperature peaks in the snapshot are also due to the adiabatic effects (Karampelas et al. 2017). Around the apex,

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4. [https://www.chiantidatabase.org/]
5. [https://wiki.esat.kuleuven.be/FoMo](https://wiki.esat.kuleuven.be/FoMo)
mixing of different regions plays an important role in equilibrating the temperature, while near the footpoint, where mixing is less strong, the high-temperature regions are more localized. The radiative cooling results in significantly lower temperatures both at the apex and near the footpoint. This is because the radiative cooling is proportional to the square of the density, and thus has most effect on the high-density loop interior. However, despite the strong radiative cooling, the radiative simulation maintains a temperature above the initial temperature of 1 MK. For the background corona, the gradual increase in temperature around the apex is due to the ponderomotive force (Terradas & Ofman 2004), which is also found in previous numerical simulations (Magyar & Van Doorsselaere 2016; Guo et al. 2019b; Karampelas et al. 2019a).

### 3.2. Volume-averaged Energy Density and Temperature

In order to investigate the evolution of global energy and temperature, we examine the volume-averaged energy density and temperature in the subdomain of $|x|, |y| \leq 4$ Mm for both simulations. The variations of internal, kinetic, magnetic, and total energy densities with respect to their initial values ($E(0)$) are shown in the top panel of Figure 4. In the adiabatic simulation, the kinetic energy first increases as the loop is oscillating as a standing kink mode. When KHI has fully deformed the loop cross section by $t \approx 500$ s and small-scale structures are formed, the kinetic energy becomes saturated. The numerical resistivity and viscosity increase the internal energy and heat the plasma by dissipating the magnetic and kinetic energy on the grid scales. By inspecting the red solid line in the top panel of Figure 4, we find that the evolution of internal energy has two stages. When $t < 650$ s, the rate of increase of internal energy is growing, while after $t > 650$ s, an almost constant rate of increase is seen. This difference is caused by the different dynamic behavior of the loop. The first stage is during the growth of KHI, with an accelerated energy transport from large to small scales. During the second stage, when KHI has saturated and the loop cross section is fully deformed and small-scale structures are generated, a constant rate of energy transfer is found as indicated by the straight red solid line.

In the radiative simulation, when radiative cooling is turned on at $t = 600$ s, the kinetic energy is not obviously influenced, while the magnetic and internal energies are changed compared to the adiabatic ones. As the gas pressure of the loop is decreased by
radiative cooling, the inward Poynting flux at the side boundaries of the simulation increases the magnetic energy of the loop. Due to the cooling, the internal energy stops increasing and stabilizes around a fixed value. Indeed, in this stage of our configuration, the driver energy input from the footpoint is approximately balancing the radiative energy losses. To better gauge the heating in the simulation, we examine the volume-averaged temperature in the subdomain, which is shown in the bottom panel of Figure 4. In the adiabatic case, the temperature increases continuously due to persistent viscous and resistive heating at the small scales. It also follows the two-stage process that we discussed before, with first a steady increase until the KHI saturates. In the latter stage, the temperature increase in the adiabatic case is nearly constant. For the radiative simulation, when radiative cooling is switched on, the temperature remains almost constant in the second, radiatively cooling stage of the simulations. This result demonstrates that the energy dissipation of small-scale structures can balance the radiative cooling and maintain a high-temperature loop. It also indicates that transverse oscillation dissipation is a promising mechanism for heating coronal loops.

To further address the heating effects, it is important to show the evolution of the internal energy of the regions where radiative cooling is stronger than background heating. To do this, we select the regions with $\rho \geq 1.1 \rho_e$ and show in Figure 5 the results of (a) volume-averaged temperature, (b) volume-averaged internal energy density, (c) total volume, and (d) total internal energy of this domain. We find that the temperature within this domain shows a similar trend to that in Figure 4, for both adiabatic and radiative cases. This result shows that the heating effects are promising to balance the radiative cooling of the denser coronal loop. However, the internal energy density decreases obviously, even in the adiabatic case, which seems to contradict the increase in temperature. A similar situation was already discussed by Magyar & Van Doorsselaere (2016) and Karampelas et al. (2017). The decrease in the internal energy density is not caused by the energy loss in this domain, but by the expansion of this domain, as Figure 5(c) shows, due to the development of KHI and the generated vortices. The gradual increase in internal energy in Figure 5(d) is also predominantly caused by the expansion of the domain.

3.3. Resistive and Viscous Heating Rates

As mentioned in previous sections, energy dissipation and heating in our model are caused by numerical dissipation. Numerical resistivity or viscosity can be estimated by conducting a group of runs including varied physical resistivity or viscosity. In this section, we estimate the numerical resistivity and viscosity in a different way by taking into account the energy equations.

The change in volume-averaged magnetic energy density $B^2/8\pi$ can be expressed as (e.g., Schnack 2009)

$$
\frac{1}{V_0} \int B^2 dV = -\frac{1}{V_0} \int \frac{e}{4\pi} (E \times B) \cdot dS dt
- \frac{1}{V_0} \int \frac{e c}{4\pi} J \cdot EdV dt
- \frac{1}{V_0} \int \frac{e c}{4\pi} [(E_0 + \eta n J) \times B] \cdot dS dt
- \frac{1}{V_0} \int (J \cdot E_0 + \eta n J^2) dV dt,
$$

where we let $E = E_0 + \eta n J$, $E_0 = -(\nu \times B)/c$ is the convective electric field and $\eta_n$ is the numerical resistivity. For simplicity,
we define the volume-averaged magnetic energy density as

\[ E_m = \frac{1}{V_0} \int \frac{B^2}{8\pi} dV. \]  

Equation (4) shows that the change in \( E_m \) is caused by the terms from ideal MHD and the terms containing the numerical resistivity \( \eta_n \). We define the change in the magnetic energy density caused by ideal MHD as

\[ E_{m,\text{ideal}} = -\frac{1}{V_0} \int \frac{c}{4\pi} (E_0 \times B) \cdot dS dt \]

\[ - \frac{1}{V_0} \int \int J \cdot E_0 dV dt. \]

Assuming \( \eta_n \) is a constant, \( \eta_n \) can thus be derived by comparing \( E_m \) and its ideal MHD counterpart \( E_{m,\text{ideal}} \). The top panels of Figure 6 show the results for the adiabatic case. We find that initially \( E_m \) matches \( E_{m,\text{ideal}} \), but later an obvious difference is seen between the two curves. This is because numerical resistivity will play a role dissipating magnetic energy when small-scale structures are generated. We see that \( \eta_n \) remains almost constant after \( t = 500 \) s. The averaged numerical resistivity in the range \( 500 \) s \( \leq t \leq 2000 \) s is \( \eta_n = 2.5 \times 10^{-9} \) s. The corresponding magnetic Reynolds number is \( R_m = 4\pi v_c l_c / (c^2 \eta_n) = 2.1 \times 10^5 \), taking the characteristic velocity as the initial internal Alfvén speed \( v_c = v_{Alf} = 3778 \) km s\(^{-1}\), and the characteristic length as the loop radius \( l_c = R = 1 \) Mm.

Figure 3. From top to bottom: plane cut of density, temperature, and \( z \)-components of current and vorticity at \( z = 0.1L \) and \( z = 0.5L \) at \( t = 860 \) s for adiabatic and radiative simulations. The animated images have the same layout as the static figure, and run from \( t = 0 \)–\( 2000 \) s.

(An animation of this figure is available.)
The change in volume-averaged kinetic energy density \( \rho v^2 \) is

\[
\frac{1}{V_0} \int \frac{1}{2} \rho v^2 dV = - \frac{1}{V_0} \iint \left( \frac{1}{2} \rho v^2 \vec{T} + \vec{P} \right) \cdot \vec{v} \, dS 
+ \frac{1}{V_0} \iiint (p \nabla \cdot \vec{v} + \vec{J} \cdot \vec{E}_0 - \overline{\Pi} : \nabla \vec{v}) dV dt 
\]

where \( \vec{P} = p \vec{T} - \overline{\Pi} \) is the total stress tensor, and \( \overline{\Pi} \) is the viscous stress tensor. For simplicity, we adopt the expression for the viscous stress tensor from fluid dynamics (e.g., Batchelor 2000),

\[
\overline{\Pi}_{ij} = \mu_n \left( \frac{\partial u_i}{\partial x_j} + \frac{\partial u_j}{\partial x_i} - \frac{2}{3} \delta_{ij} \nabla \cdot \vec{v} \right). 
\]

Here we use \( \mu_n \) to denote the numerical dynamic viscosity. Using a similar method to that above to estimate \( \eta_n \), and assuming \( \mu_n \) is also a constant, \( \mu_n \) can be derived by comparing the kinetic energy density \( E_k \) (left side of Equation (7)) and its ideal MHD counterpart \( E_{k,\text{ideal}} \) (right side of Equation (7) without the \( \overline{\Pi} \) terms). The bottom panels of Figure 6 show the results. We find that the kinetic energy density saturates at around \( t = 500 \) s, while its ideal MHD counterpart keeps increasing. The discrepancy between the two curves is caused by numerical viscosity. The averaged numerical viscosity in the range \( 500 \leq t \leq 2000 \) s is \( \mu_n = 1.1 \times 10^{-4} \) g (cm s)\(^{-1}\). The corresponding Reynolds number is \( R = \nu v/(\mu_n/\rho_0) = 1.7 \times 10^5 \), taking \( \rho_0 = 3 \times 10^8 m_p \) cm\(^{-3}\) as the density of the loop interior in the initial state.

Using the numerical resistivity \( \eta_n \) and numerical dynamic viscosity \( \mu_n \) from Figure 6, we can quantify the dissipation or heating rates caused by numerical resistivity \( D_{\text{res}} = \eta_n J^2 \) and numerical viscosity \( D_{\text{vis}} = \overline{\Pi} : \nabla \vec{v} \). Figure 7(a) shows \( D_{\text{res}} \) and \( D_{\text{vis}} \), together with the background heating rate \( H_0 \), and the
Figure 5. (a) Temperature, (b) internal energy density, (c) volume, and (d) total internal energy of the domain with $\rho \geq 1.1\rho_0$ for adiabatic (solid lines) and radiative (dashed lines) cases.

Figure 6. Top: evolution of magnetic energy density ($E_m$) and the magnetic energy density from ideal MHD ($E_{m,\text{ideal}}$), with estimated numerical resistivity $\eta_n$. Bottom: evolution of kinetic energy density ($E_k$) and the kinetic energy density from ideal MHD ($E_{k,\text{ideal}}$), with estimated numerical dynamic viscosity $\mu_n$. The results are for the adiabatic case in the fixed volume of $|x|, |y| \leq 4$ Mm.
The radiative cooling rate, averaged over the domain of \(|x|, |y| \leq 4 \text{ Mm}\) for the radiative case. Both the background heating and the radiative cooling are smoothly turned on at \(t = 600 \text{ s}\). From Figure 7(a) we see that the cooling rate and background heating rate \(H_0\) are several times larger than both \(D_{\text{vis}}\) and \(D_{\text{res}}\). This is mainly because both resistive and viscous dissipations are localized around the loop while Figure 7(a) is averaged over a domain much larger than the loop size. \(H_0\) is constant, as we mentioned above. The radiative cooling rate is roughly balanced by \(D_{\text{vis}}, D_{\text{res}},\) and \(H_0\). The radiative cooling rate decreases only slightly, meaning that the development of small-scale structures does not obviously influence the radiative loss of this domain. Figure 7(b) shows the same quantities but averaged over the domain \(\rho \geq 1.1 \rho_0\). This domain is not fixed, but rather expands as small-scale structures are generated (see Figure 5(c)). The volume expansion obviously decreases the volume-averaged radiative cooling rate. The radiative cooling rate is also roughly balanced by \(D_{\text{vis}}, D_{\text{res}},\) and \(H_0\). Even though \(H_0\) contributes around 30% of the total heating rate in the time range \(700 \text{ s} \leq t \leq 2000 \text{ s}\), the dissipations from resistivity and viscosity play the major role in the dynamics and deserve further studies in future.

4. Discussion and Conclusion

It is necessary to discuss the thermal conduction, which is not included in our simulations. Thermal conduction can also cause energy loss of a loop. In our initial setup, the timescales of a loop cooled down by thermal conduction \((\tau_c)\) and by radiative cooling \((\tau_r)\) can be roughly estimated as (e.g., Cargill 1994)

\[
\tau_c = 4 \times 10^{-10} \frac{n L_0^2}{T^{5/2}} = 1.2 \times 10^4 \text{ s}
\]

\[
\tau_r = \frac{2 n k_b T / (\gamma - 1)}{n^2 \Lambda(T)} = 5300 \text{ s}
\]

where \(n = 3 \times 10^8 \text{ cm}^{-3}\) is the loop density, \(L_0 = 100 \text{ Mm}\) is the loop half-length, \(T = 1 \text{ MK}\), and \(\Lambda(T) = 2.6 \times 10^{22} \text{ erg cm}^3 \text{ s}^{-1}\) is the radiative cooling rate at 1 MK. We have \(\tau_r > \tau_c\) in our setup, indicating that radiative cooling is more important than thermal conduction, though further analysis is required.

Thermal conduction can transport heat along the magnetic field, which may change the temperature profiles we obtained in this work. Figure 4 shows that heating can globally balance cooling, while Figure 3 shows different temperature profiles around the footpoint and the apex. In Figure 8 we show the temperatures averaged over five subdomains with the same volume but centered at different \(z\) positions for the radiative case. We find that heating is profound around the apex, while around the footpoint the temperature decreases because of cooling. Around \(z = 0.3L\) and \(z = 0.7L\), heating and cooling almost balance each other. If thermal conduction is included, the temperature profile along the \(z\) direction could be smoothed. In the adiabatic stage of \(t < 600 \text{ s}\), we obtain a similar temperature profile to Guo et al. (2019b) in their monolithic case. The small reduction in temperature near the footpoint is probably due to the ponderomotive force that can transport mass from the footpoint to the apex and thus influence the temperature profiles (Terradas & Ofman 2004).

Thus, we have proven with our simulations that AC heating mechanisms can be used in 3D MHD models of coronal loops to heat the coronal plasma and compensate for the massive radiative losses. This was done by comparing an adiabatic simulation with a simulation in which radiative cooling was switched on. In the latter simulation, the temperature was kept at coronal values, because of the balance between the dissipation of the input driver flux, the background heating, and the radiative cooling.
This model constitutes a proof-of-principle for the AC heating mechanisms that were proposed half a century ago. Our model opens the door for further exploration of these heating mechanisms. This includes models incorporating a chromosphere and transition region, which will induce field-aligned mass flows in response to the turbulent heating and thermal conduction. Moreover, we can look forward to models encompassing a whole active region heated with MHD waves. More importantly, the models that will build upon our model can be contrasted to earlier DC heating models for the solar corona. This comparison could in principle indicate which observable effects can distinguish the two heating classes in observations.

Before closing, some remarks are necessary for future studies. First, our current simulations have energy outflows near the boundary, which could be improved in the future, and second, a parametric study where the wave period is changed near the boundary, which could be improved in the future, and

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