Overdense microwave plasma heating in the CNT stellarator

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Abstract
Overdense plasmas have been attained with 2.45 GHz microwave heating in the low-field, low-aspect-ratio CNT stellarator. Densities higher than four times the ordinary (O) mode cutoff density were measured with 8 kW of power injected in the O-mode and, alternatively, with 6.5 kW in the extraordinary (X) mode. The temperature profiles peak at the plasma edge. This was ascribed to collisional damping of the X-mode at the upper hybrid resonant layer. The X-mode reaches that location by tunneling, mode-conversions or after polarization-scrambling reflections off the wall and in-vessel coils, regardless of the initial launch being in O- or X-mode. This interpretation was confirmed by full-wave numerical simulations. Also, as the CNT plasma is not completely ionized at these low microwave power levels, electron density was shown to increase with power. A dependence on magnetic field strength was also observed, for O-mode launch.

Keywords: stellarator, microwave heating, collisional damping, mode conversion, upper hybrid, finite difference time domain, overdense plasma

(Some figures may appear in colour only in the online journal)

1. Introduction

The need for high triple product and desire for high plasma β in fusion experiments translates into a simultaneous requirement for high densities and temperatures. In devices with relatively low magnetic fields, this implies that the ratio \( \omega_{pe}/\omega_{ce} \) of the electron plasma frequency to the electron cyclotron frequency might exceed unity. In this case, the plasma may not admit the propagation of electromagnetic waves in the electron cyclotron frequency range. Thus, if electron cyclotron heating [1] is desired, overdense heating techniques become necessary.

One common overdense heating mechanism involves mode-conversion of slow extraordinary (SX) electromagnetic waves to electrostatic electron Bernstein waves (EBWs), whose propagation is immune from upper density limits [2]. The conversion occurs at the upper hybrid resonance (UHR), which is reached by three main techniques: high-field side launch of an SX mode (SX-B scheme [3–5]), low-field side launch of a fast extraordinary mode (FX-B scheme, also known as FX-(SX)-B [6, 7]), or low-field side launch, with a special oblique angle, of an ordinary mode (O-X-B scheme [8–12]).

However, the excitation of EBWs requires a finite electron Larmor radius. When the temperature at the UHR is too low, the SX mode spends an infinite time at the UHR without ever departing from it and converting to EBWs [13]. Instead, it is collisionally damped [14]. This is sometimes considered an undesirable effect, in cases in which the excitation and core deposition of EBWs is preferred [2]. On the other hand, collisional damping permits to heat the overdense edge of fusion plasmas, as experimentally observed in NSTX [15].

Densities, temperatures, and thus collisionalities similar to the edge of fusion plasmas are encountered in smaller, university-scale experiments. An example is the TJ-K torus, also overdense, and also successfully heated by collisional damping of microwaves at the UHR [16]. The vacuum wavelength \( \lambda_0 \) adopted in TJ-K was longer than in NSTX, due to the lower field and hence lower \( \omega_{ce} \). At the same time, the plasma minor radius \( a \) was smaller, and the experiment had to be interpreted by full-wave modeling. For various physics reasons, full wave modeling in the electron cyclotron
frequency range is also becoming increasingly popular in fusion devices of larger $a$ and smaller or much smaller $\lambda_0$, where it is computationally more demanding [17, 18]. Here it is argued that modeling small, low-field plasmas is a natural first step on the way to simulating larger ones.

The low field, low aspect ratio CNT stellarator [19] can be viewed as a further extension of the TJ-K long-wavelength regime. CNT was originally dedicated to non-neutral and quasi-neutral plasma research [20, 21], but has recently been re-purposed to investigate error fields [22], high-$\beta$ stability [23] and image inversion [24] in neutral stellarator plasmas. As part of this, it was equipped with microwave heating at 2.45 GHz. The corresponding vacuum wavelength, $\lambda_0 = 12.2 \text{ cm}$, is comparable with the plasma minor radius $a \approx 13 \text{ cm}$ and about one-third of the major radius $R \approx 30 \text{ cm}$. On this scale, the broad launched microwave beam strikes a broad region of the plasma at a range of incident angles and polarizations. This variation has been limited to some extent in CNT by placing a waveguide very close to the plasma edge. However, no focusing element was used in this initial study. Also, the launch system was not yet optimized for any particular overdense heating mechanism.

In spite of this, plasmas were attained in CNT that were overdense to O-mode propagation by factors of more than 4, as shown in the present paper. Section 2 describes the heating and diagnostic systems employed for this work. Profiles of density and temperature confirmed overdense heating for both O- and X-mode launch. These profiles are presented in section 3 alongside a study of their dependence on power and magnetic field. Section 4 describes the full-wave numerical method used to interpret the experimental results. Interpretations are discussed in section 5, coming to the conclusion that, similar to TJ-K, the CNT plasma is heated by collisional damping of the X-mode at the UHR. This is true even when the initial launch is in O-mode, primarily due to polarization-scrambling reflections by the inner walls and in-vessel coils.

EBW heating of the CNT core will require a hotter edge, for efficient mode conversion.

2. Experimental setup

2.1. Heating system

The microwave heating system is shown schematically in figure 1(a). The microwave source is a 10 kW, 2.45 GHz magnetron manufactured by Muegge. At full power, its output is CW. Reduced power (in a time-averaged sense) is obtained by pulsed operation at roughly 5 kHz: the lower the duty-cycle, the lower the averaged output power. The magnetron launcher excites the $\text{TE}_{10}$ mode (figure 1(b)) in a rectangular waveguide.

A three-stub tuner controls the percentage of power coupled to the plasma, and the reflected power is absorbed in a water-cooled isolator upstream of the tuner (not shown). The power injected in the vessel is determined by the difference of the forward and reflected power as measured by a pair of Schottky diodes fixed to a dual directional coupler.

This net injected power is an upper bound for the power actually deposited in the plasma. The rest is dissipated on the resistive walls or leaks out of the vessel through few unshielded ports, after multiple reflections off the walls.

Downstream of the tuner is a twistable, flexible rectangular waveguide followed by a rectangular-to-circular taper that can rotate freely on the flange on its circular side. The polarization is linear in the $\text{TE}_{10}$ mode in the rectangular waveguide (figure 1(b)). This is tapered into the $\text{TE}_{11}$ mode in the circular waveguide (figure 1(c)), which is also linearly polarized, to some approximation. By varying the orientation of the twistable, flexible waveguide and rotatable taper, one can adjust the polarization angle of the injected linear polarization relative to the magnetic field in the plasma. This allows selecting the O- or X-mode as the dominant polarization injected in the plasma. During nominal O-mode (or X-mode) launch, about 80% of the power injected is actually polarized in O-mode (or X-mode). The remaining 20% mismatch is due to the injected polarization being linear but the actual O and X eigenmodes in the plasma being elliptically polarized, for oblique injection, and easily calculated from the Appleton–Hartree dispersion.

Following the taper is the launch antenna: a section of circular waveguide that leads into the vacuum vessel up to near the plasma edge. The launch antenna is effectively a re-entrant port, held at atmospheric pressure to avoid unwanted breakdown and plasma formation within the waveguide, at locations where the wave-frequency equals high-order EC harmonics. Such plasma would partly absorb or fully reflect the high-power microwaves in the waveguide, before they reach the stellarator plasma. Both effects are undesired. A quartz window at the end of the launch antenna functions as the vacuum break.

A schematic of the orientation of the launch antenna relative to the plasma is shown in figure 2. The 2.45 GHz
2.2. Diagnostics

The primary plasma diagnostic used in this work was an array of Langmuir probes similar to what was used to diagnose non-neutral plasmas in CNT [20]. The probe tips are made of halogen light bulbs with the glass removed to expose the tungsten filaments. The tips are mounted on a ceramic rod which can be moved longitudinally into and out of the plasma with an edge-welded bellows drive actuated by a recently installed stepping motor, at a speed of \( \sim 1 \text{ cm s}^{-1} \). The probe enters the plasma in a wide cross-section and therefore must move about 30 cm to scan from the edge to the axis (the average minor radius is 13 cm). For comparison, plasma discharges may last up to 45 s, limited by the heating of the coils, although more typical discharges of only 8–10 s are examined here. The probe array intersects the plasma between \( \phi = 180^\circ \) and \( \phi = 215^\circ \), far from the launch antenna which aims at \( \phi = 90^\circ \).

Electron temperature and density were determined from probe current–voltage characteristics \( I(V) \), obtained by sweeping the probe bias with a repetition rate of 200 Hz (that is, every 5 ms). The steady-state measurements reported in this paper were derived from averaging 0.25–0.5 s of data.

The effective minor radius (flux coordinate) of the probe at each longitudinal position was determined using an electron beam/phosphor rod technique commonly used to visualize flux surfaces [22, 26, 27]. In this case, the probe itself emitted the electron beam, and the flux surface images were aligned with previous measurements, thereby associating each probe position with a three-dimensional flux surface. The geometry of these surfaces is well understood following a diagnosis of CNT’s error fields [22].

Electron temperature \( T_e \) and density \( n_e \) are assumed uniform on the flux surfaces. Outside the last closed flux surface (LCFS), however, few measurements are available and the extrapolation of \( n_e \) to locations close to the launch window is subject to uncertainties. Models of magnetized plasma sheaths near solid surfaces predict that a magnetic pre-sheath will extend into the plasma by a distance given by the sound speed divided by the ion cyclotron frequency, \( c_s/\omega_{ci} \) [28]. Near the window, this distance is of order 1 cm for singly ionized Ar, which is the working gas used in these experiments. The layer in which the density drops has little effect on wave-propagation, due to its thinness compared to the wavelength and other scales of interest. Additionally, that layer is highly underdense to O-mode.

3. Profile measurements: evidence of overdense heating

Microwave plasmas were generated with the heating system described in section 2.1 and diagnosed with the probes of section 2.2, finding evidence of overdense plasma heating.

Electron temperature and density profiles were obtained for various heating powers and magnetic field strengths. Each profile was obtained from a number of discharges realized with the same heating and backfill pressure parameters, while
the probe was scanned through the plasma and scrape-off layer. In a typical 8–10 s long discharge, the power gradually increased as the magnetron warmed up, reaching its flat top at \( t = 4 \) s. However, some plasma parameters were observed to evolve from \( t = 5 \) s onwards, especially in high-power discharges. Therefore, profile measurements were restricted to the interval \( t = 4–5 \) s in each discharge. To generate profiles, the probe was radially scanned by \( 1 \) cm during that interval, as well as from one discharge to the other. Tests confirmed the discharge-reproducibility to be well within the \( n_e \) and \( T_e \) error bars.

3.1. Dependence on heating power

3.1.1. O-mode launch. Figures 3(a)–(b) show \( n_e \) and \( T_e \) profiles for argon (Ar) plasmas heated with O-mode waves, for different values of injected power \( P \). As discussed in section 2.1, this is an upper bound for the power absorbed in the plasma. The profiles were fitted with 7th-order polynomials with the constraint of \( \frac{dn_e}{dx} = \frac{dT_e}{dx} = 0 \) at \( x = 0 \), for reasons of symmetry.

The backfill pressure was \( (1.4 \pm 0.2) \times 10^{-5} \) Torr in each case. Consequently, the mean-free-path of neutrals before being ionized was \( \gtrsim 10 \) cm. As the plasma minor radius was \( a \approx 13 \) cm, the neutral density was assumed to uniformly evaluate \( n_n = 4.5 \times 10^{17} \) m\(^{-3} \) throughout the plasma, corresponding to the backfill pressure.

For the most part, in the cold plasmas considered the Ar atoms were either singly ionized or not ionized at all, and there were very few ions of higher charge. Hence, by quasi-neutrality the ion density was \( n_i \approx n_n \). In turn, \( n_e \) plotted in figure 3(b) was comparable with the neutral density just

![Figure 3. Results for O-mode launch at various power levels. (a)–(b) Langmuir probe measurements of \( n_e \) and \( T_e \) projected on the \( x \) axis (as defined in figure 2), fitted with 7th-order polynomials. The vertical gray line denotes the LCFS. The horizontal dashed line is the cutoff density for O-mode. The solid black curve represents a typical \( n_e \) profile from CNT’s earlier 1 kW heating system, with N\(_2\) instead of Ar as the working gas. (c) \( dW/dV \) data and spline fits. (d)–(e) Contours of cutoffs and resonances in the \( xy \) plane (figure 2) for 0.5 and 8.0 kW heating power. Regions overdense to the O-mode are shaded; portions that could be accessible by O-X conversion are shaded lighter. The solid, dashed and dotted green lines indicate the axis and two contours of relative intensity (at \( 1/e \) and \( 1/e^2 \) of the on-axis values) of the microwave-beam, as it would propagate in vacuum. The star in (d) is a reference for the discussion in section 5.2. The vacuum wavelength is shown in (e). (f)–(g) Contours of time-averaged electric field and percentage of first-pass power absorption determined by the full-wave code for 0.5 and 8.0 kW.](image-url)
estimated. Consequently, \( n_1 \approx n_{\text{ne}} \), i.e. the gas was only partly ionized. As a consequence, increasing the heating power did not result solely in heating (higher \( T_e \)), but also in more ionization (higher \( n_e \)), as seen in figures 3(a)–(b). Therefore, as a metric of effectiveness of power-coupling to the plasma, it is convenient to also plot, in figure 3(c),

\[
\frac{dW}{dV} = \frac{3}{2} n_1 k_\text{B} T_1 + \frac{3}{2} n_e k_\text{B} T_e + n_1 E_{\text{ioniz}} \approx n_1 (1.95 k_\text{B} T_e + E_{\text{ioniz}}).
\]

Here \( k_\text{B} \) is the Boltzmann constant and \( E_{\text{ioniz}} \) is the first ionization potential. For Ar it is \( E_{\text{ioniz}} = 15.8 \) eV, clearly not negligible in the energy balance of the partly ionized plasmas of \( k_\text{B} T_e = 2–7 \) eV presented here. The quantity \( dW/dV \) is the volumetric density of energy \( W \) stored in the plasma, expressed in terms of thermal energy of the ions, of the electrons, and energy expended in the ionization process. The latter is the stored energy associated with the steady-state number of ions per unit volume, \( n_e \), regardless of their temperature. Ionization and recombination rates are not considered because, in steady-state, their contributions to the power balance cancel out. The assumption \( T_e \approx 0.3 T_i \) was made for the ion temperature in the electron-heated CNT plasma, and discussed in a previous paper [23].

The profiles of \( T_e \) and \( dW/dV \) (figures 3(b)–(c)) peaked at the edge, and peaks were found to depend on the heating power \( P \), suggesting power-deposition at the edge.

Density profiles (figure 3(a)) show that nearly the entire plasma was overdense to O-mode propagation, at all power levels. The region of \( T_e \) peaking and likely power deposition was also overdense.

Also shown in figure 3(a) is the \( n_e \) profile from an earlier nitrogen plasma produced with a rudimentary heating system, consisting of a 1 kW magnetron placed in front of a viewport. Compared to that, \( n_e \) in the new experiments increased by over an order-of-magnitude, even when using as little as 0.5 kW (in blue in figure 3(a)). This indicates that the new launcher, closer to the plasma and better aimed, deposits a much higher fraction of the injected power into the plasma. A further improvement could be brought by a focusing, steerable mirror. Focusing is beneficial in general, and steering of a collimated beam is beneficial for the O-X-B scheme. As for the temperature, \( T_e \) also increased with the new launcher, but less significantly (figure 3(b)). This is due to the plasma not being fully ionized, as just discussed. Therefore, it is not surprising that increasing the power coupled to the plasma primarily results in more ionization, i.e. higher \( n_e \).

Figures 3(d)–(e) show contours of cutoffs and resonances in the \( xy \) cross-section defined in figure 2. The contours were calculated from the \( n_e \) profiles for O-mode launch at the lowest and highest power levels. Also shown are the launch window and the propagation axis and width of the ‘vacuum microwave beam’ (i.e., the microwave beam as it would look in the absence of plasma). As indicated by the O-mode cutoff curves, both cross-sections are almost entirely overdense—thus, evanescent—to the O-mode. However, some O-X conversion may occur and make the lighter-shaded regions partly accessible to the SX mode, up to a ‘turning point’ (not shown) where it bends toward the UHR.

Finally, shown in figures 3(f)–(g) are full-wave simulations realized with the IPF-FDMC code [29]. These simulations will actually be discussed in sections 4–5, but it is convenient to plot them here, next to the experimental profiles (figures 3(a)–(c)) and accessibility plots (figures 3(d)–(e)) which they refer to.

### 3.1.2. X-mode launch

Figure 4 presents the results of another power scan, for the same Ar gas pressure, but launching an X-mode. The shapes and trends in the \( n_e, T_e \) and \( dW/dV \) profiles are similar to the O-mode case. Again, heating seems to occur predominantly at the plasma edge. The low, medium and high power levels were obtained using the same stub tunings as in the O-mode power scan; however, the values of \( P \) were different. This indicates that X-mode coupling to the plasma differed from O-mode coupling. This was expected, due to the different reactivity of the plasma to the two modes, which, in turn, is related to the different locations of the O- and X-mode cutoffs. Cutoffs and evanescent regions for the two modes are plotted respectively in figures 3(d)–(e) and 4(d)–(e).

In generating those figures, \( n_e \) and \( T_e \) were extrapolated outside the LCFS as discussed in section 2.2. As a result, contours outside the LCFS (based on radial extrapolations) are less accurate than inside the LCFS, which are based on Langmuir probe measurements, although at a different toroidal location, far from the launcher (figure 2(b)). For O-mode launch, this inaccuracy should not impact the analysis because \( n_e \) is usually well below cutoff outside the LCFS, hence refraction of the O-mode is negligible and insensitive to relatively small \( n_e \) variations. The effect on X-mode propagation might be more significant because \( n_e \) outside the LCFS is on the order of the FX cutoff density. Hence, an error in density could switch the medium from overdense to underdense for the X-mode. Yet, even if the region in front of the launcher becomes evanescent for the X-mode, as in the case of figure 4(d) (as opposed to figure 4(e)), finite tunneling is still possible, provided the region is not too thick.

It should also be noted that, for simplicity, the cutoff and resonant layers in figures 4(d)–(g) were determined assuming propagation at 90° relative to the magnetic field. This is not the case everywhere, due to the incident beam being angularly broad (due to the low frequency) and the curvature of \( B \) being appreciable over the transverse size of the beam (due to CNT’s low aspect ratio and to the beam being broad). As a consequence, in different locations the propagation vector \( k \) forms a different angle with \( B \), and different from 90°. The evanescent region between the UHR and FX cutoff varies accordingly (figure 5) and becomes thicker for shallower incidence. This, combined with the larger distance traveled across the layer for more oblique angles, reduces the X-mode tunneling efficiency for grazing incidence.
3.1.3 Dependence of global parameters on heating power

The density and temperature profiles in figures 3–4 increased with the injected heating power, but maintained their shapes nearly unaltered. Therefore, their dependence on power can be summarized by only plotting the volume-averaged density \( n_e \) and edge temperature \( T_{e,90} \), as in figures 6(a)–(b). Here \( \langle n_e \rangle = (1/V) \int n_e(\rho) \, dV \), where \( n_e(\rho) \) is the fitted experimental profile, which is a function of the flux surface coordinate \( \rho \). The temperature \( T_{e,90} \) is the fitted \( T_e \) evaluated at 90\% of the effective minor radius. That was the approximate location where \( T_e \) tended to reach its maximum. Hence, \( T_{e,90} \) is effectively the peak temperature, or close to the peak temperature.

The data points in figures 6(a)–(b) were all obtained using \( |B| = 88 \) mT. They are a combination of all data from figures 3–4 as well as the 88 mT data from a \( |B| \) scan to be presented in figures 7–8 and discussed in section 3.2.

The temperature increases with the heating power \( P \), as expected (figure 6(b)). The density also increases with \( P \), approximately like \( P^{0.27} \) (figure 6(a)). This \( \langle n_e \rangle \) increase is due to the plasma not being fully ionized, as discussed in section 3.1.1.

The energy stored in the plasma and the volume-averaged plasma beta were calculated as

\[
W_{\text{plasma}} = \int n_e(\rho)[1.95k_BT_e(\rho) + \epsilon_{\text{ioniz}}] \, dV, \tag{2}
\]
was the same as for the power scans given in the legends. A scan had the values are the highest.

Changes of and shows that the power actually deposited in the plasma can be estimated as (for the coupled power equation). As expected, both quantities increase with heating power and, therefore, values of , establishing a more favorable power balance, resulting in significantly higher temperatures and, therefore, values of [23].

Finally, the energy confinement time (figure 6(e)) was estimated as . Note that is an upper bound for the coupled power (section 2.1): numerical modeling shows that the power actually deposited in the plasma can be 20%–85% lower (tables 1–2). Therefore, the plotted is really a lower bound for the confinement time. In addition, assumptions and experimental errors in equation (2) propagate into systematic errors and uncertainties in (equation (2)) and (equation (3)) and .

Energy confinement is found to decrease with like (figure 6(e)). As mentioned above, power-balance (hence, energy-confinement) in the partly ionized plasmas considered here suffers from severe radiative losses. One driver of radiative losses, electron–ion recombination, increases roughly in proportion with and thus in proportion with in a quasi-neutrality plasma. Hence, it is not surprising that decreases (figure 6(e)) as increases (figure 6(a)). Indirectly, this could also explain why decreases with (due to the correlation between and ).

Future power upgrades are expected to yield temperatures in excess of the radiation barrier, at which radiative losses will diminish. As a result, is expected to increase dramatically compared to figure 6(e) and not to follow the rough scalings just discussed, but rather the ISS04 international scaling [31], as discussed in [23].

3.2. Dependence on magnetic field

A series of experiments was also conducted in which the heating power was kept constant but the strength of the magnetic field, , was varied. The purpose of this was to move the cyclotron resonance and, to a lesser extent, the UHR and the X-mode cutoffs, but not the O-mode cutoff. These changes affected the efficiency of candidate mode-conversion and heating mechanisms. Therefore, the scan had the promise of helping to isolate which mechanisms were taking place in the experiments.

The intermediate was the same as for the power scans in section 3.1. The higher and lower values are the highest and lowest permissible values such that neither the fundamental nor the second harmonic resonant surfaces intersect the launch window (similar to figure 2(a)). This is to avoid damage to the window. The values of given in the legends of figures 7–8 are evaluated at the origin of the xy plane.

The stub-tuning was the same for all the experiments in this scan. This resulted in net power coupling of 1.0 ± 0.2 kW in all cases except X-mode launch at 88 mT, which had a net coupling of about 2 kW.

The results are shown in figure 7 for O-mode launch and in figure 8 for X-mode launch. For O-mode launch, the and profiles exhibit the greatest sensitivity to . Changes of profile shape were also observed, with both becoming increasingly hollow as increased, and the core density being halved.

By contrast, for X-mode launch the profiles exhibited little to no variation with .

4. Full-wave modeling

In order to predict and interpret the experimental results, the interactions between injected microwaves and the CNT plasma was modeled by means of the full-wave, finite-difference time-domain code IPF-FDMC [29]. The code solves Maxwell’s equations coupled with the fluid equation of motion for the electrons in a nonuniform magnetized plasma. The unknowns are the wave electric and magnetic field and the electron current density , as functions of two spatial dimensions (2D) and time. As a result, IPF-FDMC accounts for effects such as the O-X mode conversion and tunneling of the X-mode through the evanescent region between the UHR and fast X cutoff. Since the code makes no
assumption about the scale length of the plasma relative to the wavelength, it is well-suited to simulate the propagation of long waves in the compact CNT plasma, and their reflection by the CNT walls and in-vessel coils. Here the 2D simulations are sufficient to determine the heating mechanism and qualitatively understand the role played by reflections. 3D full-wave simulations would be quantitatively more accurate, but they are notoriously demanding from a computational standpoint, and thus very rare [32], especially in the electron cyclotron range. They are left as future work.

IPF-FDMC was extensively benchmarked against cold plasma theory [33], other full-wave codes [34, 35] and a wave-kinetic equation solver [36]. In addition, IPF-FDMC was benchmarked against measurements in WEGA [11] and TJ-K [16]. Predictions were also made for TJ-II [29], RFX-mod [37] and Pegasus [38], and should be compared with measurements when these will become available.

A damping term invoking a collision frequency $\nu$ in the range $10^{-5} < \nu/\omega < 10^{-3}$ is assumed, which damps the slow X-mode in the vicinity of the UHR [16]. Here $\omega$ is the wave frequency. The total damped power is not sensitive to the value of $\nu$ within this range. As will be discussed in section 5.3, collisional damping is expected to dominate over conversion to EBWs, in the CNT experiments presented here. No cyclotron damping is accounted for in these simulations, but is expected to be negligible anyway (section 5.2).

Full-wave calculations were performed for all values of $P$, $|B|$ and polarizations in the experimental scans of section 3. Selected cases (maximum and minimum $P$ and $|B|$), corresponding to panels (d) and (e) of figures 3–4 and 7–8, are plotted respectively in panels (f) and (g) of the same figures. Shown are the contours of steady-state, time-averaged amplitude of the microwave electric field, $|E|$, normalized to its peak value (which changes from one contour plot to the other). The profiles of $n_e$ and $T_e$ as functions of $x$ and $y$ used in the simulations are based on the fits of the experimental profiles in panels (a) and (b).

The launcher is located on the right of the computational domain and injects linearly polarized waves with a Gaussian distribution of intensity. This is a realistic model of the experiment (figures 1(c) and 2(a)). Perfectly absorbing conditions are adopted at all other boundaries, unless noted otherwise. Hence, the quantity $|\langle E \rangle|$ plotted so far refers to the first pass of microwaves but neglects additional power re-impinging on the plasma after reflections off the vessel walls or internal coils. The

![Figure 7. Similar to figure 3, but for the scan of magnetic field strength in the case of O-mode launch. Each plasma coupled to 1 kW of heating power.](image-url)
percentages of first-pass power absorption, primarily due to collisional damping, are indicated in panels (f) and (g).

In each case in which the UHR is visible, enhancement of $\mathcal{E}$ is noticeable in a narrow region (much narrower than the vacuum wavelength) close to the UHR. Only the X-mode is sensitive to this resonance, but enhancement is observed for nominal O-mode launch as well. This is ascribed to a combination of O-X mode conversion and lack of modal purity: as discussed above, neither in the experiment nor in the modeling are the linear polarizations injected pure O- or X-polarizations, which would be elliptical. This implies that some X-mode is injected during nominal O-mode launch.

Field enhancement near the UHR favors collisional damping. In some cases this can account for as much as 36% absorption in a single pass (figure 8(f)). In other cases, only 1% of power is absorbed at the first transit (figures 3(g) and 7(f)). This is due to the recessed location of the UHR, such that, in those cases, the beam does not encounter the UHR in its first transit.

4.1. Modeling multiple reflections

In the experiment, the unabsorbed power is reflected by the metallic walls of the vacuum vessel and by the metal cases of the in-vessel coils, and effectively re-injected in the plasma. To study this effect, the calculations of figures 7(g) and 8(g) were repeated with perfect electrical conductors lined on the left and bottom edges of the computational domain. The conductors simulate the in-vessel coils placed at the approximate same locations, although with an inclination of 78° (or, equivalently, 102°) relative to each other (figure 9).

For O-mode launch, this results in field-enhancement and absorption in a new location (upper left of figure 10(a)), compared with figure 7(g). The field pattern for X-mode launch (Figure 10(b)), on the other hand, does not change significantly with respect to the same case with fully absorbing boundaries (figure 8(g)). In brief, reflections may or may not contribute to field-enhancement and absorption, depending on geometry and polarization. The same is expected of wall reflections.

These effects were integrated in 2D full-device simulations. The computational domain is now a cross-section of the entire vessel, which the plasma intersects twice. One result in this larger domain, for the case of O-mode at low power with $B = 87.5$ mT, is shown in figure 11(a). Significant field enhancement is visible at the lower left of the figure. As there...
is no direct line of sight from the launcher to that toroidally remote location, the wave can only reach it after multiple reflections off the walls and coils.

As a result, the fraction of absorbed power increases from 2% for first-pass (figure 3(f)) to 51% when multiple reflections are taken into account (figure 11(a)). The remaining 49% of the power was reflected back into the waveguide, as evidenced by the standing-wave pattern in the launch antenna in the upper-right corner of figure 11(a). Other standing waves can also be noticed, between the plasma and the wall. All 12 experimental combinations of $P$, $B$, and polarization of figures 3–4 and 7–8 were numerically modeled. For brevity, contours of $|E|$ are only shown for eight cases in the restricted domain with absorbing boundaries (panels (f)–(g) of the said figures) and for three cases in the enlarged domain, with reflections included (figure 11). However, significant increases of fractional absorption were observed in all 12 cases (tables 1–2). Barring three exceptions at 15%, 29% and 39%, the fraction of power absorbed amounted to 50%–80%. These calculations underscore the crucial role of reflections in coupling the microwave power to the CNT plasma in these experiments.

Even higher absorption percentages are expected in 3D simulations. This is because the fractional surface of the launcher and other unshielded ports (relative to the total surface of the wall) is even smaller than the fractional arc subtended by the launcher (relative to the total circumference of the wall) in the 2D problem. Therefore, on average in 3D the beam will experience more reflections before impinging on a port and abandoning the vessel. Consequently, it will cross the plasma a higher number of times, and deposit more power.

Resistive losses at the wall are neglected because the reflectivity of stainless steel to 2.45 GHz microwaves exceeds 99.9% [39].

It is well known that an incident O-mode (or X-mode) can be partly reflected by smooth metallic surfaces as X-mode
(or O-mode). This property is known as depolarization or polarization mixing, and in the fusion literature it is often referred to as polarization scrambling [40]. When using reflective boundaries, as in figures 10 and 11, IPF-FDMC naturally accounts for this effect.

Corrugated surfaces can cause additional scrambling [41], to the point that depolarization is often used as a measure of surface roughness [42]. This is because waves of different polarizations tend to be reflected at the corrugations’ crests and troughs, respectively. Said otherwise, realistic features, unless much smaller than \( \sim \lambda/10 \), introduce phase-shifts that affect polarization. In the present case, thanks to the large wavelength \( \lambda \), it was sufficient to model the CNT vessel and coils with \( \sim 1 \) cm accuracy to take this effect into account as well.

Yet, the amount of polarization scrambling observed in the simulations was relatively modest. For instance, the X wave launched in figure 11(b) remains fairly purely polarized in X-mode even after reflections, as indicated by the fact that it barely penetrates beyond the right-handed FX cutoff.

Finally note that if single-pass absorption is high (figure 8(f)), very few reflections occur before most of the power is completely absorbed by the plasma or reflected back in the waveguide, as illustrated in figure 11(c).

5. Discussion: UHR accessibility, mode conversions and heating mechanisms

In the following subsections various mode-conversion and wave-damping mechanisms are examined one by one, and invoked or discarded as possible explanations of the experimental results presented in figures 3–8, on the basis of analytical or numerical arguments.

5.1. UHR access and mode conversions

First-pass SX-B conversion is not possible because the beam is launched from the low-field side. Yet, the X-mode can access the UHR from the high-field side after multiple reflections by the overdense plasma, the chamber walls and the in-vessel coils, which are coated in metal jackets. This is possible even for O-mode launch, because ‘polarization scrambling’ [40] at the walls and coils converts part of the injected O-mode in X-mode.

The X-mode can access the UHR from the low-field side as well. For this it has to tunnel across the evanescent layer that extends from the FX cutoff to the UHR [43]. This is a pre-requisite for the FX-(SX)-B scheme. The power transmission \( T^2 \) through this layer is approximately \( \exp(-\eta m_j) \), where \( \eta = [2\pi \Delta x/\lambda D] \) is the Budden factor (section 13.5 of [44]) and \( \Delta x \) the thickness of the layer. As an example, along the beam-axis in figure 8(e) it is \( \Delta x = 0.5 \) cm, yielding \( T^2 = 45\% \). That is, nearly half the power injected in X-mode along that ray can tunnel through the evanescent region and reach the UHR. However, in other locations the evanescent layer is thicker (for example, at higher \( y \) in figure 8(e)) and can become even thicker due to density fluctuations, resulting in some cases in \( T^2 \) as small as 0.1%.

Incidentally, for sufficiently high \( n_e \) the FX cutoff and UHR lie ‘behind’ the launcher, there is no evanescent layer to penetrate, and the beam impinges directly on the overdense core, as shown in figure 4(e).

Finally, let us discuss the X-mode accessing the UHR via O-X conversion. The injected beam is angularly broad. Hence, parts of it will not have the special direction needed for O-X conversion. This will be improved in the future by a refocusing, steerable mirror. In addition, the injected
polarization is linear, whereas a pure O-mode propagating obliquely to \( \mathbf{B} \) would be elliptically polarized, which will require the introduction of a \( \lambda/4 \) phase-shifter. Nevertheless, for low wave frequency, \( \tau_0 \) would be nearly two orders of magnitude longer. As a consequence, \( \tau_{O1} < 10^{-4} \) and \( \tau_{X1} < 10^{-7} \), hence first-pass cyclotron damping is negligible in the underdense edge of CNT.

Fusion plasmas are typically optically thick to first-harmonic O-mode, thanks to their high \( T_e \). In CNT, though, \( \nu_i^2 \) is three orders of magnitude lower. Additionally, due to the low \( |\mathbf{B}| \), hence low \( \omega_{ce} \), the low wave frequency, \( \lambda_0 \) is nearly two orders of magnitude longer. As a consequence, \( \tau_{O1} < 10^{-4} \) and \( \tau_{X1} < 10^{-7} \), hence first-pass cyclotron damping is negligible in the underdense edge of CNT.

The portions of the beam marked in light green in figure 12 propagate across the overdense core (in gray). They are obviously attenuated, due to the overdense region being evanescent but could still deposit appreciable power by cyclotron damping, provided that they reach the Doppler-broadened EC resonance with finite wave-amplitude and, again, provided that the local optical depth is sufficient. With regard to \( \tau \), the increased value of \( \omega_{pe}/\omega_{ce} \) is advantageous, but is not high enough to compensate for the five orders of magnitude mentioned before. Moreover, the wave evanesces too rapidly. As an example, let us consider the 0.5 kW O-mode case in figure 3, where \( n_e \approx 1 \times 10^{12} \text{ m}^{-3} \), and let us consider a location where \( |\mathbf{B}| \approx 90 \text{ mT} \) and \( \theta \approx 90^\circ \). The Appleton–Hartree dispersion relation \( \omega_1^2 \) gives us a wave-number \( k \approx 30 \pi / \lambda_0 \) for O-mode propagation; hence, the electric field decays to \( 1/e \) of its incident value within 3 cm of the cutoff layer. The penetration depth would be even lower in denser plasmas.

This is in contradiction with O-mode experiments, where denser plasmas tend to be hotter (figures 3(a)–(b) and 7(a)–(b)). Due to this contradiction, as well as to the low optical thickness and short penetration length, it is concluded that cyclotron damping of the O- and X-mode cannot occur in significant amounts neither in the underdense nor in the overdense regions. While it may contribute a small population of non-thermal electrons, it cannot be the dominant heating mechanism.

5.3. Collisional absorption at the UHR versus excitation of EBWs

In section 5.1 we have discussed three mechanisms by which the X-mode can reach the UHR despite cutoffs. It remains to be discussed whether it converts to EBWs or, as we will conclude, it is locally absorbed due to collisions.

Note that electromagnetic EBWs have wavelength comparable with the electron Larmor radius \( r_L \), propagate away from the UHR (in this heating scheme; the opposite in a diagnostic scheme) and are\( \text{cyclo} \text{tron}-\text{damped} \) near the EC resonance.

By contrast, the electromagnetic SX mode approaches the UHR, decelerates, acquires shorter and shorter wavelength but does not convert to EBWs, because first it is collisionally damped near the UHR.

In CNT, like in other cold plasmas or plasma edges, collisional damping prevails on the conversion to EBWs. This is because the temperature is too low \( [13] \) or, equivalently, the electron–ion collisionality is too high \( \nu_{ei} = 1-10 \text{ MHz} \). This implies that collisions are numerous, during the long time spent by the slow X wave near the UHR. The time spent near the UHR can be estimated numerically, but, as an order of
magnitude, panel (b) in figures 3–4 and 7–8 indicates temperature of order $T_e = 5eV$ at the UHR location. This, combined with $|B| \approx 50 \text{ mT}$, implies electron Larmor radii of order $\rho_e = 80 \mu\text{m}$. Waves of frequency $\omega = 2.45 \text{GHz}$ and wavelength $\lambda = 5\rho_e$ propagate with phase velocity $v_p = 10^6 \text{ m s}^{-1}$. By the time this wave has traveled 10 cm near the UHR, on average every electron has collided with an ion once, in a plasma of $n_e = 10 \text{ MHz}$. Here ‘every electron’ includes electrons oscillating in the wave field (upper hybrid oscillations) and, effectively, sustaining these partly electrostatic waves. Due to collisions, energy is transferred from these electrons (thus, ultimately, from the wave) to the ions, which are not resonating with the EBW and are not supporting it. In brief, electron–ion collisions subtract energy from the wave as it approaches the UHR, depleting it of most of its energy and largely preventing mode conversion to EBWs. Here this effect is quantified with the aid of the IPF-FDMC full-wave code that, as mentioned, includes collisions and would be able to resolve the conversion from SX to short-wavelength EBWs, if this were taking place.

6. Conclusions and future work

In summary, overdense microwave plasma heating has been observed in CNT with $n_e$ exceeding the cutoff density by factors of more than 4. Density and temperature profiles tend to be hollow, and changes in heating power affect temperatures primarily at the edge rather than at the core. Variations in $|B|$ had significant effects on the profiles for O-mode launch, but not for X-mode launch.

These observations are consistent with collisional damping of an X-mode at the UHR. Note that the X-mode can reach the UHR even for O-mode launch, after O-X mode conversion and/or polarization-scrambling reflections off the CNT walls and internal coils. Such interpretation is supported by full-wave modeling performed with the IPF-FDMC code. In the relatively cold plasmas presented, heated with less than 10 kW of microwave power, the X-mode is completely damped before completing its conversion in the Bernstein mode.

Future work will include upgrades to higher power for high $\beta$ stability research [23], as well as modulated-power experiments that will improve our understanding of the heating deposition locations. In addition, full-wave calculations in three dimensions (3D) will further refine the 2D modeling presented here. 3D effects are expected to be important, due to the low aspect ratio of CNT.

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