

Helicon plasma mode transitions at low magnetic fields

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The nature of a low-field helicon plasma mode transition is examined using experiment and theory. Tests in a radio frequency plasma source reveal an abrupt transition between helicon and capacitive discharge modes that occurs at a critical magnetic field strength. To characterize the mode transition, a global model for low-field $m = 0$ helicon plasmas is developed based on power absorption via Landau damping. A general scaling law for the critical field strength is derived from the model and experimentally validated using measurements of electron density and temperature. The assumption of power absorption via Landau damping in the model leads to a critical condition for the ratio of wave phase velocity and electron thermal velocity. Measurements of wave phase velocity and electron temperature while operating on mixtures of argon and N_2O support the theoretical prediction and show that the mode transition depends on electron properties rather than on the ion and neutral composition. Agreement between theory and experiment shows that Landau damping is the primary power absorption mechanism in low-field, conducting boundary helicon plasmas.

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I. INTRODUCTION

Radio frequency (rf) coupled plasmas are used in a wide range of applications across the fields of material processing [1–3], electric propulsion [4–7], and fusion [8,9]. This versatility stems in part from the fact that rf energy can be coupled to the plasma through several distinct mechanisms, leading to different discharge modes [10]. The helicon mode (W mode) is of particular interest due to its efficient energy coupling, resulting in a high density discharge on the order of 10^{18} m^{-3} [11,12]. Furthermore, in applications including propulsion and material processing, the use of low-field helicon discharges in the range of 0–100 G [13,14] can be advantageous due to their reduced magnet power requirements [14–16]. Studies of low-field helicon plasma sources (PSs) have observed mode transitions that depend on power and axial magnetic field [17–21], and although energy absorption of high-field helicons has been extensively studied [11,22–24], uncertainty remains regarding the underlying physics of low-field helicon wave coupling. To explain the nonlinear phenomenon observed in low-field discharges, mechanisms including electron-cyclotron resonance [14,25,26], mode conversion [14,22,27], wave-particle trapping [25,28], and end-wall wave reflection [16,25,29] have been proposed. In this article, we explore how energy coupling in low-field helicons relates to an observed mode transition using experiment and theory.

II. RF PLASMA SOURCE

The low-field helicon mode was investigated using the plasma source shown in Fig. 1(a). The PS is made of 304 stainless steel, is 28 cm in length, and has an internal diameter of 10 cm. An electromagnet provides an axial magnetic field, creating a magnetic nozzle (MN) [30–33]. The field strength at the center of the glass window (B_0) is 16 times weaker than that at the MN throat. Note that B_0 is the magnetic field minimum within the discharge region. Driving the rf discharge is a two-turn spiral antenna powered by a 13.56 MHz rf generator and automatic matching network; this antenna geometry has been found to excite the $m = 0$ helicon mode [13]. Strong impedance matching was obtained for all operating conditions (<5% reflected power). Gas is injected through eight radial ports and the discharge exhausts into a cylindrical dielectric chamber measuring 0.8 m across and 1.3 m in length. The PS was run on mixtures of argon and nitrous oxide (N_2O) gas with a constant chamber pressure of 0.8 mTorr, commensurate with other experiments [6,22,24,28]. An rf-compensated Langmuir probe (LP) [34] was used to measure electron temperature and density [35]. Additionally, a one-dimensional inductive field probe array (B -dot probe) consisting of 12 coils arranged parallel to the PS axis was used to measure azimuthal magnetic field oscillations [12,13,36]. This configuration enabled measurements of the helicon wave structure and absorption rate.

A helicon discharge was observed in the PS, as evident by direct measurements of wave propagation using the B -dot array, density measurements consistent with other helicon discharges [37], and visual characteristics such as intensity and color [12]. A low-density, capacitive mode (E-mode) discharge was also observed [10]. For a given power and gas mixture, the E-mode plasma could be sustained across a wide range of magnetic field strengths. However, an upper limit to the magnetic field was found for the helicon discharge. As B_0

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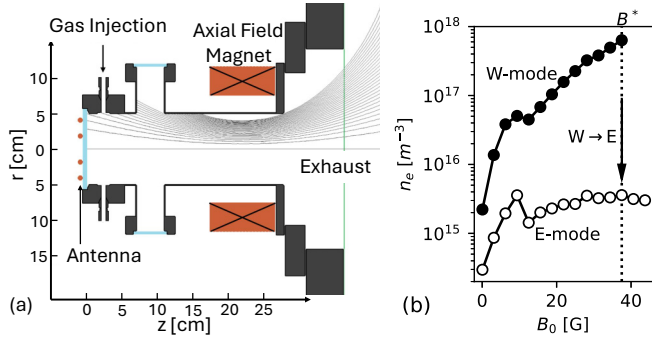


FIG. 1. (a) Geometry of experimental rf PS. (b) Experimentally measured density vs magnetic field strength. The mode transition of interest occurs as B_0 (field at center of glass window) increases above B^* , highlighted by the arrow. Data obtained for PS operating on argon with 500 W of applied rf power.

increases past a critical value, an abrupt W-E mode transition occurs. The critical field strength for the transition (denoted by B^*) depends on power and gas mixture but is repeatable for fixed conditions. Figure 1(b) shows the electron density of a 500 W argon plasma for both modes as a function of B_0 . We see that for $B_0 > B^*$ only the capacitive mode exists. In contrast to the W-E mode transition, the transition from E mode to W mode is not induced in a consistent manner. Note that although inductive coupling plays an important role in similar low-field devices [38], we did not observe a transition to the inductive mode in our experiment. When the W-E mode transition occurs, wave propagation ceases and n_e drops by over 2 orders of magnitude. This is in contrast with the response of high-field helicons where mode transitions induced by an increasing magnetic field correlate with increasing electron densities [12]. The abrupt W-E mode transition is also distinct from the previously studied low-field density peaks [19,25,39,40] and we will refer to it as the low-field mode transition.

Values of B^* versus rf power are presented in Fig. 2(a) for various gas mixtures. We see that B^* increases monotonically with both argon mass flow percentage (at a fixed power) and rf power (for a fixed gas mixture). Electron density (n_e^*) and temperature (T_e^*) near B^* for the W mode are shown in Figs. 2(b) and 2(c), respectively. Similar to B^* , n_e^* and T_e^* increase with rf

power. At low powers, a more severe drop in temperature and density is observed for N₂O-rich plasmas (0% and 25% argon mass flow) compared to B^* . While B^* and n_e^* both increase monotonically with the argon mass flow percentage, T_e^* has a nonlinear response. Figure 2(c) shows that at high rf powers, the introduction of N₂O leads to an increase in T_e^* ; however, this trend reverses for N₂O-rich plasmas.

To understand the trends in B^* , n_e^* , and T_e^* , we consider power loss mechanisms. The five test conditions in Fig. 2(b) with the lowest density also had relatively low B^* (~ 10 G), leading to elevated cross-field diffusion [1,38]. Diffusion plays a critical role in helicons [33] and represents a key power loss mechanism, explaining the drop in n_e^* . Furthermore, the presence of N₂O introduces additional reactions such as electron impact dissociation [41]. The inelastic nature of these reactions reduces the total power available to generate and heat the plasma [4,42]. Indeed, we see n_e^* and B^* increase with rf power but decrease with N₂O flow fraction. However, the nonlinear response of T_e^* shows that the response to chemistry effects is not simple. Due to the complexity of these various power loss mechanisms, we look to theory to determine the relationship between B^* and the electron properties.

III. MODE TRANSITION MODEL

The observations above suggest that the mode transition may be driven by nonlinear power coupling between the antenna and plasma. To test this theory, we adopt the power balance method of Chabert and Braithwaite [10], which considers the global balance between rf power absorption and various power losses. Loss mechanisms include ionization, radiation, wall losses, and beam losses through MN acceleration. Because these mechanisms scale primarily with density [10], we use the approximation $P_{\text{loss}} = c_{\text{loss}} n_e$, where P_{loss} is the total power leaving the system and c_{loss} is an unknown constant that represents power lost per unit density. We model the total absorbed rf power as the sum of capacitive (P_{abs}^E) and wave (P_{abs}^W) coupling: $P_{\text{abs}} = P_{\text{abs}}^E + P_{\text{abs}}^W$ (inductive coupling is ignored for simplicity). An empirically driven model for capacitive absorption [10] is assumed: $P_{\text{abs}}^E = c_E/n_e$, where c_E is a constant. Finally, for steady-state discharges, power balance is required: $P_{\text{abs}} = P_{\text{loss}}$.

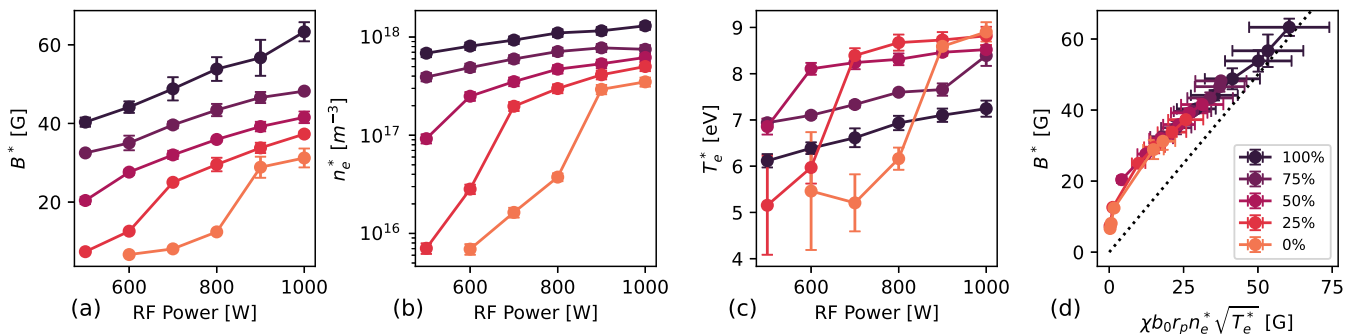


FIG. 2. (a) Mode transition critical magnetic field strength, (b) electron density, and (c) electron temperature vs power for W-mode plasma near the mode transition. (d) Comparison of the experimentally measured B^* to theoretical prediction. Data shown for different gas mixtures with the legend indicating argon mass flow percentage.

To determine P_{abs}^W , we consider the dispersion relation of a cylindrically bounded helicon with wave frequencies between the ion and electron cyclotron frequencies ($\omega_{\text{ci}} \ll \omega \ll \omega_{\text{ce}}$) [10]:

$$kk_z = \frac{\omega\omega_{\text{pe}}^2}{c^2\omega_{\text{ce}}} \left(1 - \frac{i\nu_{\text{eff}}k}{\omega_{\text{ce}}k_z} \right). \quad (1)$$

Here, k , k_z , ω_{pe} , and c represent the wave number, axial wave number, plasma frequency, and speed of light, respectively. Note that for simplicity, we ignore the electron inertia correction to the dispersion relation [14] and instead follow the approach taken by Lafleur *et al.* [25]. For an $m=0$ helicon plasma, the radial wave number satisfies $k_r = \alpha_1/r_p$ [11], where r_p is the plasma radius and $\alpha_1 \approx 3.83$ is the first zero of the Bessel function. Here, ν_{eff} is the effective collision frequency of the wave. Under the assumption that power absorption is dominated by Landau damping as initially proposed by Chen [11], we take $\nu_{\text{eff}} = \nu_{\text{ei}} + \nu_w$, where $\nu_{\text{ei}} \approx 2.9 \times 10^{-11} n_e T_e^{-3/2}$ is the electron-ion collision frequency and $\nu_w = 2\sqrt{\pi}\xi^3\omega \exp(-\xi^2)$ is the collisionless Landau damping frequency. The quantity ξ is proportional to the ratio of wave phase velocity (v_ϕ) to electron thermal velocity ($v_e \equiv \sqrt{eT_e/m_e}$): $\xi \equiv \sqrt{2}v_\phi/v_e$. Wave power absorption in a plasma of length L then becomes $P_{\text{abs}}^W = P_{\text{rf}}[1 - \exp(2L/\alpha_z)]$, where k_z is found from Eq. (1), $\alpha_z = \text{Im}(k_z)^{-1}$ is the characteristic absorption length, and P_{rf} is the applied rf power. In addition to the equilibrium condition ($P_{\text{abs}} = P_{\text{loss}}$), a stable solution to this global model also requires $dP_{\text{abs}}/dn_e < dP_{\text{loss}}/dn_e$.

The power balance model results are shown in Fig. 3 for conditions that approximate our PS. Figure 3(a) shows the dependence of ν_{eff} , ν_{ei} , and ν_w on n_e . We note the strong nonlinearity in ν_{eff} due to the contribution of ν_w . The influence of B on the shape of ν_{eff} versus n_e is shown in Fig. 3(b). Here, increasing B pushes the peak in ν_{eff} to higher values of n_e . This shift occurs because the shape of ν_w depends only on $\xi \sim v_\phi/v_e \sim B/(n_e T_e^{1/2})$. The influence of increasing B on P_{abs} is presented in Fig. 3(c). Here, we again observe a strong nonlinearity in P_{abs} versus n_e that is driven by ν_w and shifts toward higher n_e with increasing B . The role of nonlinear power absorption in the observed mode transition is illustrated by plotting P_{loss} versus n_e for a fixed value of c_{loss} . We see in Fig. 3(c) that two stable equilibria exist for low magnetic fields: W mode and E mode. However, as B increases, the shift in P_{abs} toward higher values of n_e eliminates the W-mode solution, leaving only the E mode as a stable solution. Notably, the strong nonlinearity in P_{abs} versus n_e predicts an abrupt transition to the E mode, which agrees with our experimental observations. Finally, Fig. 3(d) shows how the equilibrium values of n_e depend on B . Here, we highlight the strong similarities with the experimental measurements shown in Fig. 1(b), despite the simplifying assumptions of the model.

We seek a generalized solution to the power balance model to correlate measurements of B^* to n_e^* and T_e^* . It is apparent from Figs. 3(b) and 3(c) that the leading edge of the P_{abs} curve corresponding to the transition point correlates with the point of maximum collisionless Landau damping. The damping frequency, ν_w , has a well-defined maximum

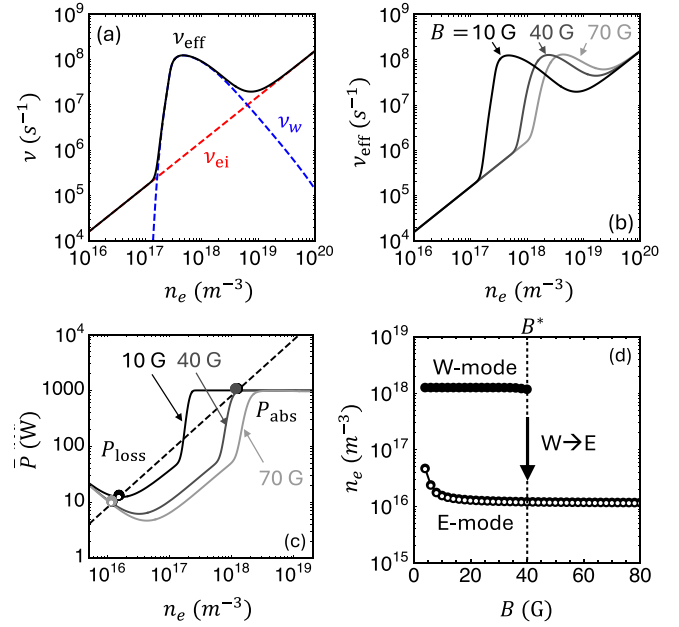


FIG. 3. Results from the theoretical power balance model: (a) collision frequencies vs density; (b) effective collision frequency vs density for different magnetic field strengths; (c) power absorption and power loss vs density for different magnetic field strengths; and (d) density vs magnetic field strength, with equilibrium W-mode (solid dot) and E-mode (open dot) solutions. The arrow in panel (d) highlights the W-E mode transition that occurs with increasing B . Results represent the following model parameters: $r_p = 5$ cm, $L = 28$ cm, $P_{\text{rf}} = 1$ kW, $T_{\text{ev}} = 7$ eV, $f = 13.56$ MHz, $\omega = 2\pi f$, and $c_{\text{loss}} = 8 \times 10^{-16} \text{ W m}^3$.

at $\xi_M = \sqrt{3/2}$; we denote the axial wave number satisfying $\xi_M = \sqrt{2}v_\phi/v_e$ as $k_{z,M}$. Furthermore, Eq. (1) can be approximated as $k_{z,M} \approx (\omega\omega_{\text{pe}}^2)/(c^2\omega_{\text{ce},M}k_r)$ provided $\nu_{\text{eff}} \ll \omega_{\text{ce}}$ and $k_z \ll k_r$. Solving for the field strength gives $B_M = b_0 r_p n_e T_e^{1/2}$, where $b_0 = (\sqrt{3}e^{3/2})/(2\alpha_1\epsilon_0 c^2 m_e^{1/2})$ is a constant.

Note that the transition does not correspond precisely to the peak in ν_w but rather to a field strength slightly greater than B_M . To characterize this shift, we solve the model over a range of conditions to quantify the ratio $\chi \equiv B^*/B_M$, which we refer to as the shape factor of P_{abs} . The transition field strength and its corresponding density (n_e^*) depend on c_{loss} ; however, estimates based on diffusion, MN approximations, and ion properties are of questionable accuracy. Instead, we can solve for B^* and n_e^* implicitly using the conditions that both power balance is satisfied ($P_{\text{abs}}(n_e^*) = c_{\text{loss}} n_e^*$) and P_{abs} is tangent to P_{loss} at the transition ($P'_{\text{abs}}(n_e^*) = P'_{\text{loss}}(n_e^*) = c_{\text{loss}}$, where $P'_{\text{abs}}(n_e^*)$ and $P'_{\text{loss}}(n_e^*)$ are the derivatives of P_{abs} and P_{loss} with respect to n_e). By substitution, we arrive at $P_{\text{abs}}(n_e^*) - n_e^* P'_{\text{abs}}(n_e^*) = 0$. Solving for n_e^* over a range of parameters creates a database of $[B^*, n_e^*, T_e^*, L, P_{\text{rf}}]$ sets that are self-consistent at the mode transition.

Solving B_M for each $[n_e^*, T_e^*]$ pair in the database reveals the shape factor to be roughly constant: $\chi \approx 1.9 \pm 0.4$. Uncertainty in χ was determined by sensitivity analysis over each parameter of the model (B^* , n_e^* , P_{rf} , T_e^* , and L). Using the shape factor, we find that the general scaling law for the mode

transition's critical field strength is given by

$$B^* = \chi b_0 r_p n_e^* \sqrt{T_e^*}. \quad (2)$$

It follows that $k_z^*/k_{z,M} = \chi^{-1}$, where k_z^* is the axial wave number at the mode transition. This suggests that the mode transition occurs at a fixed ratio of the phase velocity and electron thermal velocity: $v_\phi^*/v_e^* = \sqrt{3}\chi/2$. Here, v_ϕ^* and v_e^* denote the phase and electron thermal velocities at the mode transition, respectively. We note that according to the dispersion relation, v_ϕ increases with B . Within the helicon mode, $v_\phi/v_e \leq \sqrt{3}\chi/2$ and the velocity ratio increases with B until the critical value is reached and the transition occurs. Although identified as more accurate for low-field helicons, we note that the inclusion of electron mass in the dispersion relation is a higher-order correction that does not significantly alter the derivation of Eq. (2) or the relationship between velocities given above.

IV. EXPERIMENTAL VALIDATION

Figure 2(d) shows the measurements of B^* versus the predictions from Eq. (2). A line of unity slope is plotted for reference, representing perfect agreement with Eq. (2). We see that over the range of test conditions, the data collapse onto a single curve that roughly agrees with theory. As discussed above, the use of molecular gasses introduces loss mechanisms that change the plasma parameters and B^* . Critically, molecular dissociation leads to uncertainty in the ion composition of the plasma as new species may be present, including N, O, N₂, and NO. The fact that measurements of B^* collapse onto a well-defined curve indicates that the physical mechanism governing the mode transition is independent of gas composition and is instead governed by the electron parameters. We note that the experimental results are slightly shifted in comparison to theory, which may be a result of the LP being positioned at the MN throat. An axial sweep of the LP revealed higher densities near the antenna, implying that agreement with Eq. (2) is better in the upstream region. It is worth highlighting that even though the magnetic field is assumed to be uniform in the model and is nonuniform in the experiment, agreement is found. The wave absorption length scale was found to be small compared to the gradient length scale of B ; thus, a uniform magnetic field approximation is valid.

Our mode transition model is based on wave absorption via Landau damping, an assumption that has been highly debated in the literature, and one that warrants greater scrutiny. Early studies proposed that the efficient power coupling observed in helicon plasmas was explained by Landau damping of the wave [11]. However, later studies raised doubts about the role of Landau damping, instead concluding that power deposition is governed by coupling of the helicon wave to the TG mode [14,22–24], which has a higher damping rate [43,44]. In fact, despite being the first to propose wave absorption via Landau damping, Chen eventually concluded: “In any case, the hypothesis of Landau damping as originally proposed by Chen [...] is almost certainly incorrect and should now be discarded” [24]. The agreement between our theoretical model and experimental results suggests that the role of Landau damping in helicon wave absorption should be revisited.

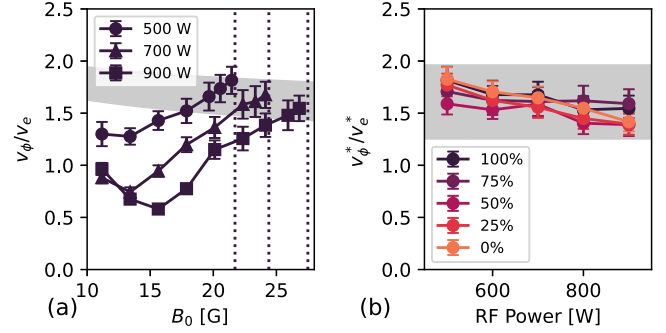


FIG. 4. (a) Ratio of phase velocity to electron thermal velocity vs applied field strength. B^* denoted by vertical dotted line for each power. (b) Ratio of phase velocity to electron thermal velocities at the transition field strength vs rf power; legend gives argon mass flow percentage. Gray regions indicate expected range of $\sqrt{3}\chi/2$ predicted by the theoretical model.

Two key predictions from our model are that (1) the mode transition occurs when v_ϕ/v_e reaches a critical value and (2) the nature of the mode transition is independent of gas species. To experimentally test these predictions, we measure v_ϕ and v_e using the B -dot array and LP, respectively. The B -dot array was positioned to measure the azimuthal wave component just downstream of the antenna. After integrating and scaling the coil signals [45], a Hilbert transform was used to provide instantaneous amplitude and phase information. The phase velocity was then calculated using the phase difference, frequency, and coil spacing; this was then time averaged to arrive at v_ϕ . Propagation of uncertainties revealed v_ϕ/v_e to have a maximum uncertainty of 10%.

Figure 4(a) shows v_ϕ/v_e versus the applied magnetic field. The vertical dashed lines denote the measured B^* for each power level, and the gray region represents the range of $\sqrt{3}\chi/2$ predicted by the model. We see in Fig. 4(a) that the velocity ratio increases with B_0 to around the critical value predicted by the model, at which point the mode transition occurs. This behavior is consistent with the physical interpretation of Landau damping, where waves couple energy into “resonant particles” that have velocities close to the phase velocity [46]. As the velocity ratio increases, the phase velocity is shifted further into the tail of the electron velocity distribution, causing a reduction in the Landau damping rate. This leads to a corresponding shift in the P_{abs} curve as seen in Fig. 3(c). When B_0 exceeds B^* , power absorption via Landau damping is insufficient to sustain the high density discharge, causing an abrupt mode transition.

Measurements of v_ϕ^*/v_e^* versus power are shown in Fig. 4(b) for various Ar/N₂O mixtures. We see that v_ϕ^*/v_e^* falls within the expected range predicted by the theoretical model (shown in gray), a result that holds for all gas mixtures tested. Additionally, measurements taken at various pressures in the range of 0.6–1 mTorr (data not presented) also showed agreement with the critical velocity ratio predicted by the model. Although the electron temperature was found to decrease with pressure, a corresponding change in v_ϕ^* occurred, resulting in a consistent velocity ratio. These results confirm that the mode transition occurs at a fixed ratio of the phase velocity and

electron thermal velocity regardless of ion and neutral species composition.

V. DISCUSSION

We have shown that the low-field mode transition occurs when the helicon wave phase velocity increases beyond a critical value, after which wave absorption via Landau damping is insufficient for sustaining a high-density discharge. This result raises the following question: Considering that Landau damping was previously dismissed as a dominant absorption mechanism in helicon experiments, why would it play a key role in our experiment? The dismissal of Landau damping as the primary absorption mechanism came largely from the observation of a strongly damped quasioleostatic wave (TG mode) that is excited by linear mode conversion of the helicon wave [14,22–24]. We note that these findings were based on experiments in which the plasma was radially bounded by an insulating wall. Shamrai and Taranov [22] used theory to show that for low-field discharges where the plasma column has a conducting boundary, such as the one considered here, the helicon and TG resonances are well separated resulting in minimal excitation of the TG mode. Thus, power absorption facilitated by coupling between the helicon and TG modes is not expected to play a role in our device.

Notably, Shamrai and Taranov refer to volumetric wave absorption via Landau damping as “problematic” due to the long wavelengths and high phase velocities observed in early helicon experiments. However, our use of an $m = 0$ antenna means that the axial wavelength is free of the minimum constraint imposed by an $m = 1$ antenna [13]. As a result, it is possible for helicon waves to propagate with $v_\phi \lesssim v_e$ at low magnetic fields, thus becoming susceptible to collisionless Landau damping. Furthermore, it is the spiral antenna geometry that enables an end-on configuration with a conducting boundary along the length of the plasma source. Therefore, our findings provide experimental confirmation for the theory of Shamrai and Taranov, and identify Landau damping of helicon waves as the primary absorption mechanism for $m = 0$ rf plasmas that are bounded by conducting walls and possess sufficiently low magnetic fields.

We note that Landau damping of TG waves in an $m = 1$ rf plasma source with insulating boundaries was previously reported by Sato *et al.* [19]. The authors showed that a density peak occurs when the wave phase velocity approaches the electron thermal velocity, suggesting enhanced power absorption via Landau damping of the TG wave. We come to a similar conclusion regarding the importance of Landau damping, but with several key differences relating to the nature of the damped wave and its influence on the global power balance within the plasma. First, we observe damping of the

helicon wave rather than the TG wave, a fundamental difference arising from the plasma boundary as argued above. Second, since the helicon wave is very weakly damped by electron-ion and electron-neutral collisions (unlike TG waves) [11,22], significant wave absorption only occurs at conditions where the discharge can support Landau damping of the wave. rf power absorption within our experiment therefore possesses a strong nonlinearity with density, which when coupled with the overall power balance gives rise to an abrupt mode transition at a critical magnetic field strength, as illustrated in Fig. 3.

VI. CONCLUSION

The results presented here represent the first experimental evidence of Landau damping being the primary power absorption mechanism for helicon waves in a plasma source. Theoretical and experimental results indicate that $m = 0$ helicon waves can propagate in plasmas with conducting radial boundaries and low applied magnetic fields. Landau damping of these waves by plasma electrons supports high power coupling efficiencies provided the magnetic field remains below a critical value associated with the low-field mode transition. The nature of the low-field mode transition is shown to be independent of the ion and neutral composition of the plasma; thus, the model and scaling law derived in this work are broadly applicable to plasma sources operating on molecular gases, such as those found in the fields of material processing and electric propulsion. Finally, we note that early investigations of Landau damping in helicon plasmas focused on finding evidence of high-energy electrons using Langmuir probes [24,47,48]. Excessive rf noise and probe heating near the antenna prevented these types of measurements in our experiment. Future work could apply noninvasive diagnostic techniques, such as incoherent Thomson scattering, to connect the wave absorption observations made here to the resulting changes to the electron energy distribution.

ACKNOWLEDGMENTS

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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