Helicons—The Past Decade

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***Abstract*— First observed in gaseous plasmas in the early 1960's, helicon discharges lay like a sleeping giant until they emerged in the 1980's, when their usefulness as efficient plasma sources for processing microelectronic circuits for the burgeoning semiconductor industry became recognized. Research on helicons spread to many countries; new, challenging, unexpected problems arose, and these have spawned solutions and novel insights into the physical mechanisms in magnetized radio-frequency discharges. Among the most baffling puzzles were the reason for the high ionization efficiency of helicon discharges and the dominance of the right-hand polarized mode over the left-hand one. The most recent results indicate that a nonobvious resolution of these problems is at hand.**

***Index Terms*—Helicon, plasma discharges, plasma processing, plasma source, plasma waves.**

I. INTRODUCTION

IN keeping with the precedent established, we begin with reminiscences of how all this came about. It was in 1962, during a sabbatical, that Chen visited Peter Thonemann at Harwell and was immediately imbued with the excitement of a new discovery: the existence of very low-frequency electromagnetic waves, called helicon waves, in a gaseous plasma. This wave, along with all the other waves, was later incorporated into the course notes for a graduate course in plasma waves and instabilities, but remained as an interesting curiosity. In retrospect, British and Australian theorists in the period following this visit did a prodigious amount of work on helicons, and the advances in the past decade have been largely due to the invention of the personal computer.

During a second sabbatical in 1985, Chen broke away from the fusion group in Canberra long enough to visit Boswell's laboratory. There, he found a small linear device with modest, subfusion caliber power supplies, creating a nearly fully ionized plasma of over 10^{13} cm^{-3} density with no axial confinement. It was immediately apparent that this was an order of magnitude better than anything he had previously seen, such as the plasma created by the nonresonant Lisitano coil. Chen carried this excitement to Japan on the way home, but only Tatsuo Shoji took the idea seriously. He encountered similar apathy back at UCLA, except from Robert Conn, co-founder of Plasma and Materials Technologies, Inc., which now markets helicon sources worldwide. Around 1989, Chen decided to leave the comfort of a strong and well-funded group in laser plasma interactions, which he had founded, in order to explore the possibilities of helicon discharges. The first years were extremely difficult, since plasma processing was practically unknown to federal funding agencies, and the semiconductor industry considered plasma science to be a worthless diversion. The UCLA group on low-temperature plasma physics had to be built from the ground up, literally. When moving the laboratory to a new building, Chen and Decker, who were later to receive the Maxwell and Ramo Prizes, respectively, laid the floor tiles themselves.

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II. BASIC DISPERSION RELATION

The basic dispersion relation for helicon waves is that for low-frequency whistlers confined to a cylinder with a coaxial dc magnetic field $\boldsymbol{B}_0$, under the assumptions that the electrons are infinitely light and the ions are infinitely heavy. In this case, the plasma current is entirely carried by the $\boldsymbol{E} \times \boldsymbol{B}$ motions of the electrons. Maxwell's equations then lead [153] to the following vector Helmholtz equation for the wave magnetic field $\boldsymbol{B} \approx \exp i(m\theta + kz - \omega t)$, where k is the axial wavenumber

$$k\nabla \times \boldsymbol{B} = k_w^2 \boldsymbol{B}, \qquad \nabla^2 \boldsymbol{B} + \alpha^2 \boldsymbol{B} = 0 \tag{1}$$

where

$$k_w^2 \equiv \omega\omega_p^2/\omega_c c^2 = \omega n_0 e\,\mu_o/B_0 \equiv \alpha k \tag{2}$$

ω_c being the electron cyclotron frequency, and ω_p the plasma frequency with density n_0. The quantity k_w will be recognized as the wavenumber for low-frequency whistler waves along $\boldsymbol{B}_0$ in free space. In a bounded system, k_w is not a free parameter but will have eigenvalues set by the permitted angles of propagation corresponding to the different radial modes. The notation used here is somewhat different from that used in earlier papers because it is more natural in view of recent advances in theory (Sections III, VIII, and IX). The solution of (1) in a cylinder of radius a is given by [153]

$$\begin{aligned} B_r &= A[(\alpha + k)J_{m-1}(Tr) + (\alpha - k)J_{m+1}(Tr)] \\ B_\theta &= iA[(\alpha + k)J_{m-1}(Tr) - (\alpha - k)J_{m+1}(Tr)] \\ B_z &= -2iATJ_m(Tr) \end{aligned} \tag{3}$$

where the J_m are Bessel functions and

$$T^2 = \alpha^2 - k^2. \tag{4}$$

The boundary conditions set the possible values of α to be approximately

$$\alpha \approx p_{mn}/a \tag{5}$$

where p_{mn} is the nth root of J_m, corresponding to the nth radial mode and the mth azimuthal mode. Most experiments are designed for the $m = \pm 1$ modes. The $m = +1$ mode is

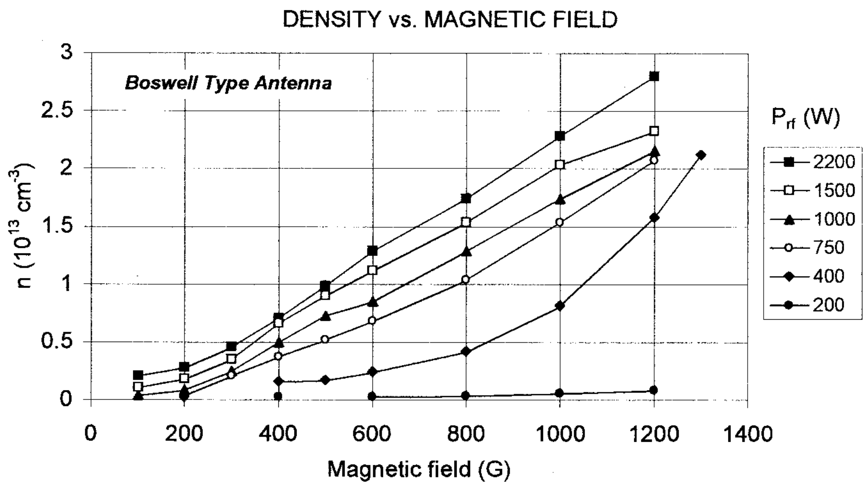


Fig. 1. Linear relationship between n_0 and B_0 [12].

right-hand (RH) circularly polarized, and the $m = -1$ mode left-hand (LH) polarized, when viewed along B_0.

According to (2), if ω and k (or the phase velocity ω/k) are fixed, the plasma density should scale linearly with the magnetic field, provided that the RF power $P_{\rm rf}$ is sufficient to produce the required density. That this basic n–B relationship is obeyed above threshold values of P_{rf} and B_0, which will be discussed further in Section VI, has been verified by a number of authors, e.g., Boswell and Porteus [6]. An example of such data is shown in Fig. 1, where a constant slope is seen above about 200 G and 400 W.

III. Nonuniform Plasmas

The formulas given in Section II are for a cylinder completely filled with a plasma of uniform n_0 and B_0. In practice, the radial density profile is more parabolic or triangular than square, and therefore the varying density had to be accounted for before detailed comparisons with experiment could be made. This was done first by Blevin and Christiansen [141], [142] and more recently by Chen *et al.* [158], [181], who included the displacement current and pointed out a numerical difficulty that it causes. Indeed, the straightforward use of the cold-plasma dielectric tensor for nonuniform plasmas is fraught with danger because singularities arise that most computer codes do not recognize.

The density profile modifies the field patterns and the eigenvalues of k_w at the density peak. A comparison of the B- and E-field patterns of the $m = +1$ and -1 modes for uniform and parabolic density profiles is shown in Fig. 2. As might be expected, the fields are more concentrated toward the axis when the density has a peak there. One sees also that the $m = -1$ mode has a narrower pattern than the $m = +1$ mode, a fact which bears on one of the most baffling features of helicon discharges which has been noticed in the last decade; namely, that the RH mode is preferentially excited under all circumstances. It is true that whistler waves, of which helicon waves are a variety, exist only in the RH version in free space; but in a confined cylinder at such low frequencies that the electron gyration direction is irrelevant, the LH mode should propagate equally well.

Although uniform magnetic fields are useful for experimental checks of theory, the fields used in helicon discharges intended for practical purposes are made to diverge away from the region of ionization in order to spread the plasma out over a wider region and to reduce the field at the target, where it may cause damage to delicate circuits. The propagation of helicon waves in a diverging field has been treated by Arnush and Peskoff [139], [140], [173]. Their method was to transform to spherical coordinates, solve in a conical section, and then join conical solutions together to match the curve of the field lines. If the plasma diverges keeping n_0/B_0 constant, the m = 0 case can be solved in closed form; otherwise, a lengthy computation is required. As the plasma radius a increases, the wavelength is observed to decrease, as required by (5), until it reaches the wavelength of a whistler wave in free space. Only rough, qualitative checks of the nonuniform-field theory have been made so far, but another interesting effect has been observed. If the antenna is located nearer one end of the discharge chamber than the other, and a "cusp" field is created in the antenna region diverging toward the short end of the machine, both the peak density and the total ionization are increased [19]. One possible explanation is that the RF energy which would have gone away from the main experimental region is somehow reflected.

IV. Wave Fields and Wave Propagation

Early measurements of helicon wave fields were made by Boswell [5] and Lehane and Thonemann [40]. Once calculations for wave fields in a nonuniform plasma became available, it was possible to compare wave profiles with theory with more accuracy. Fig. 3(a) shows magnetic probe measurements [41] of the three components of $\boldsymbol{B}(r)$ generated by a helical antenna designed to excite the $m = +1$ mode. It is seen that the points agree quite well with the theoretical curve computed for the measured density profile. When the direction of $\boldsymbol{B}_0$ was reversed, or when an antenna with the opposite helicity is used with the same $\boldsymbol{B}_0$, the $m = -1$ mode was expected to be seen. Instead, as shown in Fig. 3(b), a weaker $m = +1$ mode

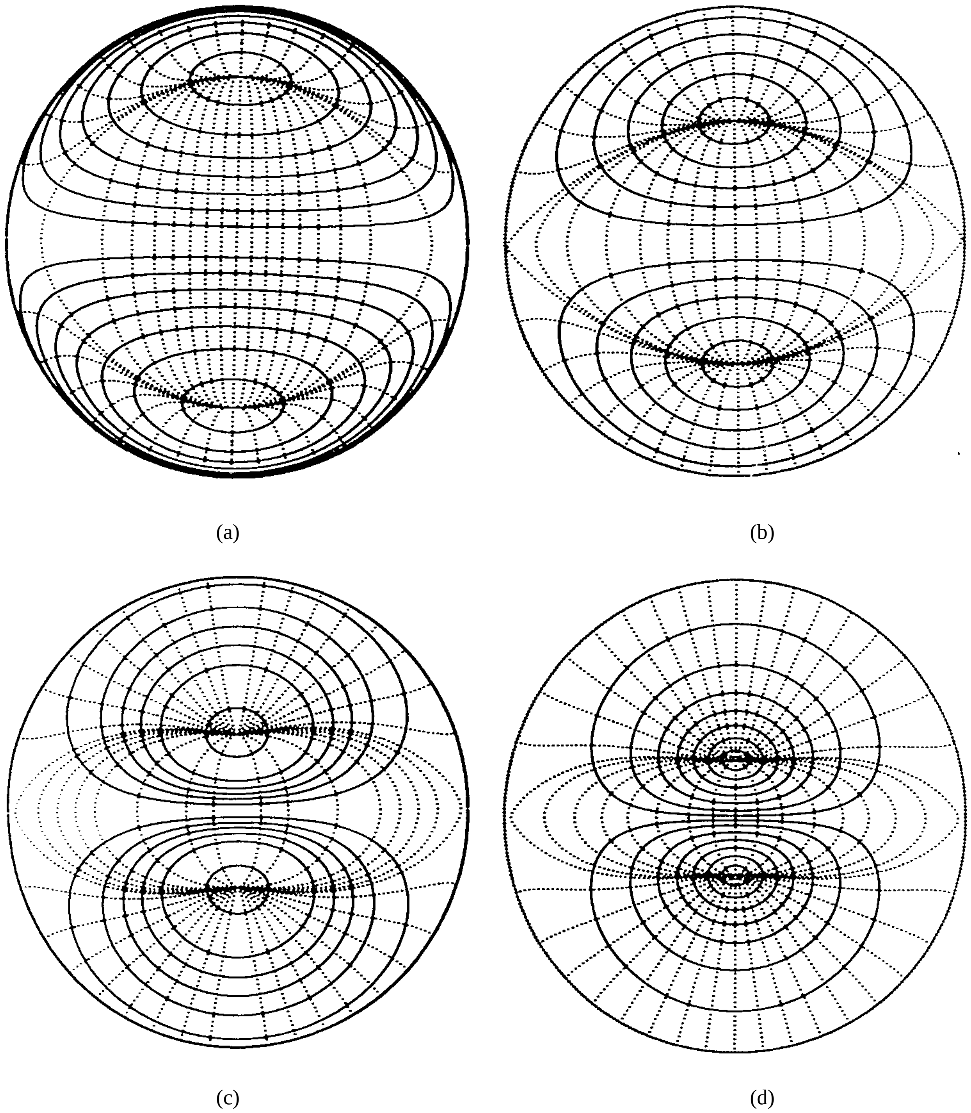


Fig. 2. Electric (dotted) and magnetic (solid) field lines for an $m = +1$ azimuthal mode (a), (b) and an $m = -1$ mode (c), (d) in a uniform plasma (a), (c) and one with a parabolic density profile (b), (d) [159].

was generated, the wave profiles differing greatly from the theory for $m = -1$. Although some authors have claimed to have seen the $m = -1$ mode [53], [54], [63], [65], [34], [49], the difficulty in exciting this mode was corroborated by other diagnostic methods and by other groups. For instance, Yasaka and Hara [65] have used bifilar antennas to produce RH and LH fields which rotate in time, and Suzuki *et al.* [63] have launched test waves with a helical antenna in a preformed plasma. A possible explanation for the predominance of the $m = +1$ mode will be discussed in Section IX. On the other hand, observations of plane-polarized patterns, which are presumably combinations of $m = +1$ and -1 modes, have been reported by several groups [51], [34]. With a plane-polarized antenna, one would expect a standing wave pattern to be generated in the azimuthal direction. The two points of view can be reconciled if $m = -1$ modes are generated in the near-field but damp rapidly away from the antenna. That this indeed happens has been confirmed by a recent experiment with rotating fields [45].

The propagation of helicon waves along $\boldsymbol{B}_0$ was first detected with magnetic probes by early workers [40], [5], and by several others in recent years [23], [26], [28], [42], [53], [63]. Boswell's measurement of $|B_z(z)|$ is shown in Fig. 4 [5]. The most obvious feature of this curve is the appearance of amplitude modulations, in this case clearly caused by standing waves. However, the modulations seen in a more recent experiment, shown in Fig. 5 [42], could not be explained as standing waves because the tube was many antenna lengths long, and the amplitude at the end was too small. Furthermore, the wavelength of the modulations did not correspond to half the wavelength of the helicon waves. Since the density varied along z (Fig. 7), this could possibly have been a case of distributed reflection, such as in the Budden problem where light entering a plasma density gradient gives rise to an Airy pattern. A simpler explanation for the pattern in Fig. 5 was available, however. If one computes the wavenumbers k of modes with various m and n numbers in a nonuniform plasma, one finds that there are two modes which lie within the peak of

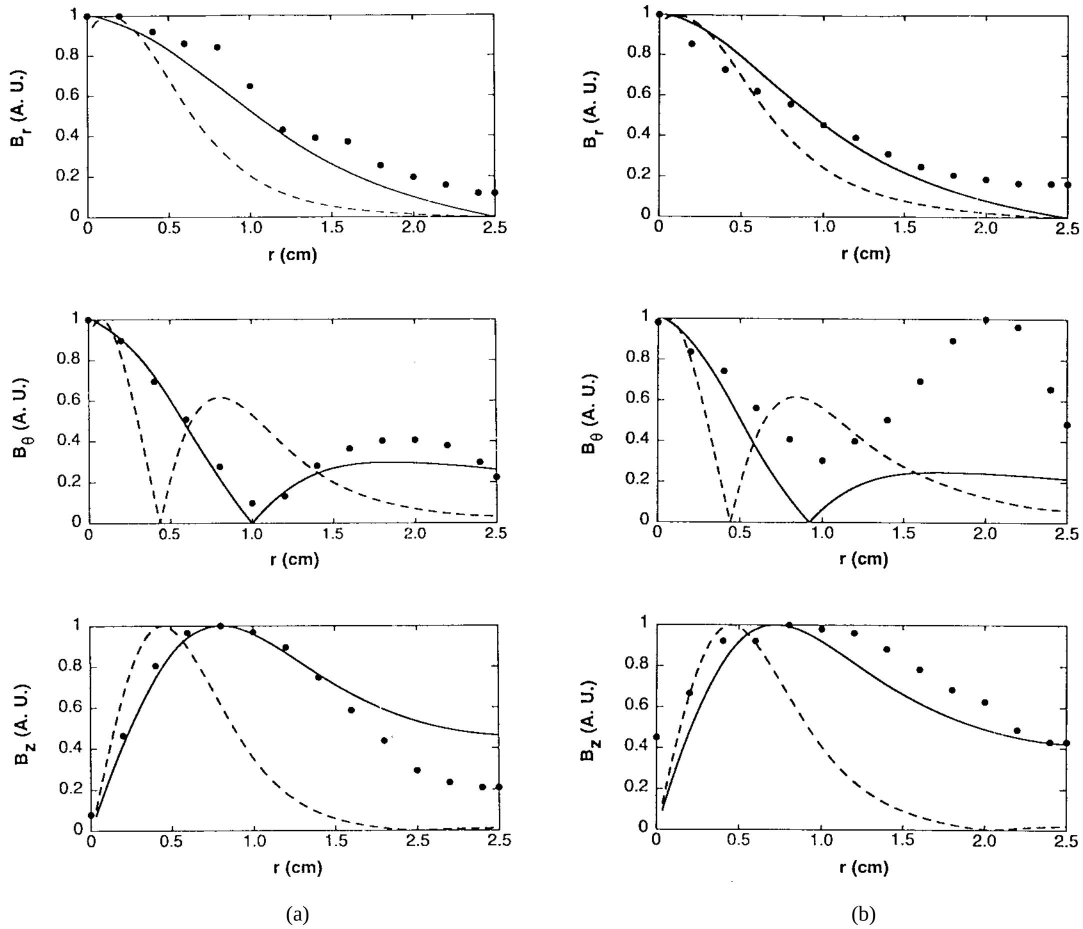


Fig. 3. Radial profiles of B_r, B_θ, and B_z as measured (dots) and calculated (lines). Antenna was set up to launch (a) $m = +1$ RH polarized waves and (b) $m = -1$ LH polarized waves. Calculations are for the $m = +1$ (solid) and $m = -1$ (dashed) modes [41].

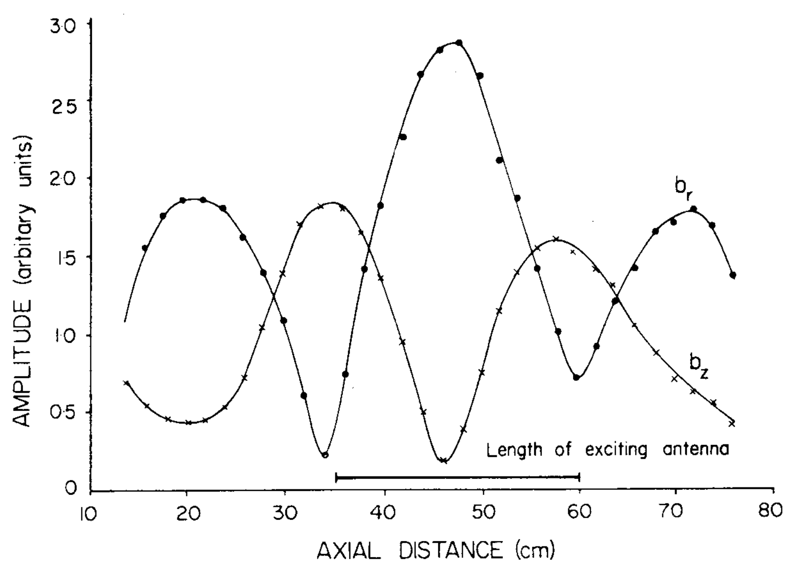


Fig. 4. Axial variation of the amplitudes $|B_r|$ and $|B_z|$ [5].

the antenna's k-spectrum. These are the first two radial modes (n = 1, 2) of the $m = +1$ RH mode, and the beat wavelength of these two modes matches that observed.

The damping length of helicon waves had been found in early work [40] to agree with theory, but this was in a weakly ionized discharge at relatively high pressure, with $B_0 \approx 150$ G, which we would now consider to be below threshold. In later work, Boswell [5] found that collision rates of order 1000 times classical would have to be invoked to explain his observations. In retrospect, this conclusion was drawn on

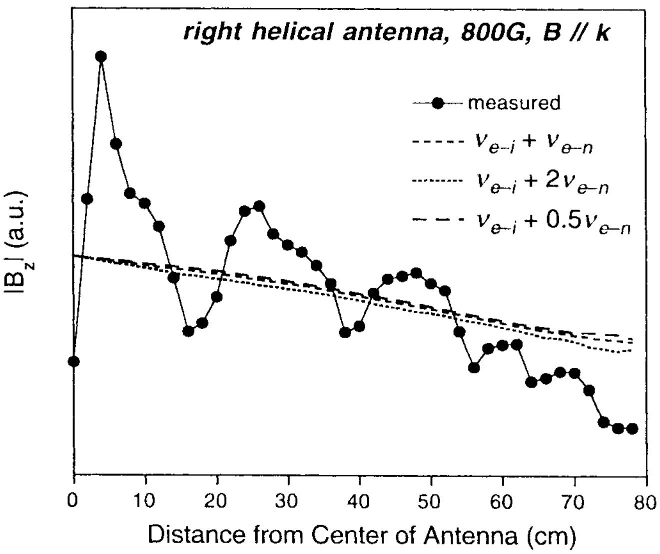


Fig. 5. Measurements of the magnitude of B_z versus distance from the antenna [42]. Also shown are curves of the theoretical damping rate calculated for electron–ion collisions plus electron–neutral collisions at various assumed pressures.

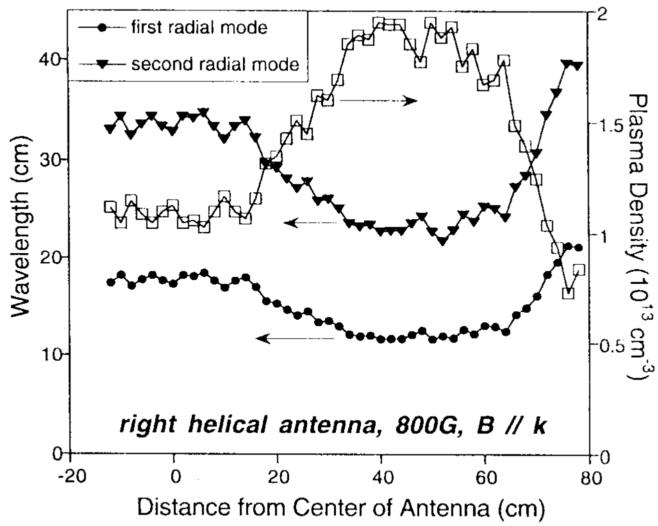


Fig. 6. Variation of the local wavelength along $\boldsymbol{B}_0$, showing anticorrelation with the local density [42].

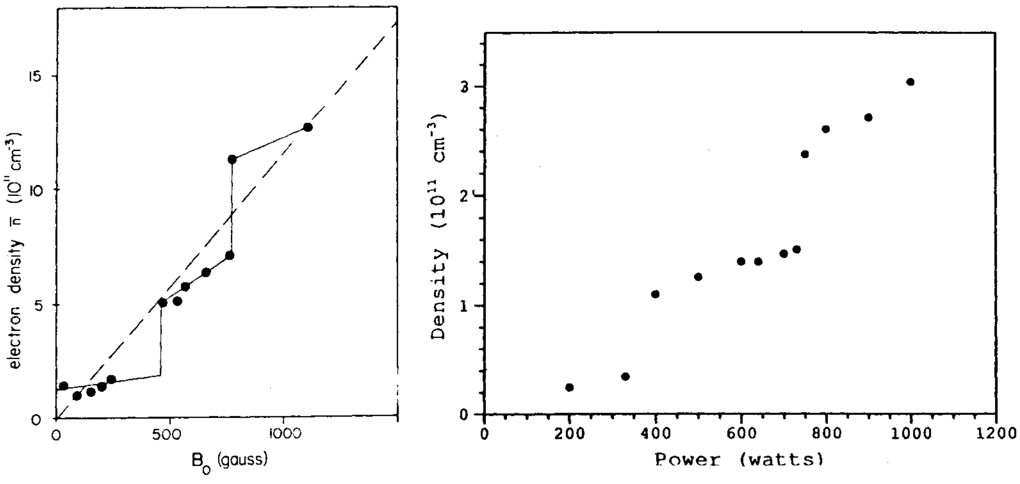
the basis of radial wave profiles, not direct measurements of the axial decay lengths. Nonetheless, it was this observation that spawned the Landau damping hypothesis discussed in the next section to explore noncollisional dissipation mechanisms. Improvements in diagnostics in recent years, particularly in the RF compensation of Langmuir probes [87], have made possible more accurate measurements of the electron temperature KT_e. Using local measurements of KT_e and n_0, the decay of wave amplitude by collisional damping could be calculated and was found to agree with measurements (when averaged over the beat-modulations), as shown in Fig. 5 [42]. In this case of high density, electron–ion collisions were dominant, as seen from the similarity of the curves when the neutral pressure was doubled or halved.

Interferometric measurements with two probes could be used to obtain the local wavelengths. Although one might expect that a dominant wavelength λ would be set by the antenna length, once the wave leaves the antenna region it is free to assume its natural wavelength. According to (2) and (5), k should be proportional to n_0/B_0 for fixed ω, so that λ should vary as B_0/n_0. At least two recent experiments bear this out. Light *et al.* [42] (Fig. 6) show λ varying inversely with n_0 in a uniform-diameter tube with a helical antenna. However, for a plasma expanding from a source into a larger chamber, Degeling *et al.* [23] have shown with a double-half-turn antenna that λ varies as $(B_0/n_0)^{1/2}$, which is expected of whistler waves in free space.

V. The Landau Damping Hypothesis

The anomalously efficient damping of helicon waves reported by Boswell [5] inspired Chen [150], [153] to suggest in 1985 that perhaps a kinetic process—namely, Landau damping—could explain the rapid transfer of RF energy into ionization events. Helicon waves have the special property that their parallel phase velocities can easily be matched to thermal velocities of electrons in the 20–200 eV range. This range is higher than typical temperatures of 2–4 eV in helicon discharges, and a straightforward computation of the Landau damping rate [153] would show that this rate is negligibly small. However, if the wave grows fast enough, it can trap thermal electrons and accelerate them to the phase velocity, as is done in RF current drive in toroidal fusion devices. These electrons would be efficient ionizers, since they have nearly the energy of the maximum in the ionization cross section. Furthermore, wave acceleration has the additional benefit that electrons can be reaccelerated after an ionization event. This idea gave rise to the design of helicon sources in which the phase velocity was matched to the optimum electron velocity for ionization [151].

Numerous attempts were made to detect the high-energy tails in the electron energy distribution function (EEDF) which would result from nonlinear Landau damping. For example, Zhu and Boswell [71] observed fast electrons by their excitation of spectral lines that have an energy threshold. Chen and Decker [11] gave indirect evidence through the charging of an endplate to high negative potentials. Loewenhardt *et al.* [43] gave clear evidence of tail electrons from Langmuir probe characteristics in a helical fusion device heated by helicon waves. Shoji *et al.* [118] showed that the EEDF broadened when ω/k was increased. R. Chen *et al.* [18] deduced from Langmuir probe data the existence of electrons whose speeds varied with the phase velocity. The authors above used probes that had inadequate or no RF compensation. At densities in the $10^{11\text{-}12}$ cm^{-3} range, Ellingboe *et al.* [24], [25] showed that the excitation of spectral lines by fast electrons occurred in pulses synchronized with the RF, as would be expected in wave-particle interactions. Such pulses were detected directly with a time-resolved gridded energy analyzer by Molvik *et al.* [46], who furthermore correlated the fast electron velocities with the helicon wave phase velocity. On the other hand, in dense plasmas above 10^{13} cm^{-3}, Blackwell and Chen [2], using carefully compensated probes, could see no deviation from a pure Maxwellian distribution. This null result may have been



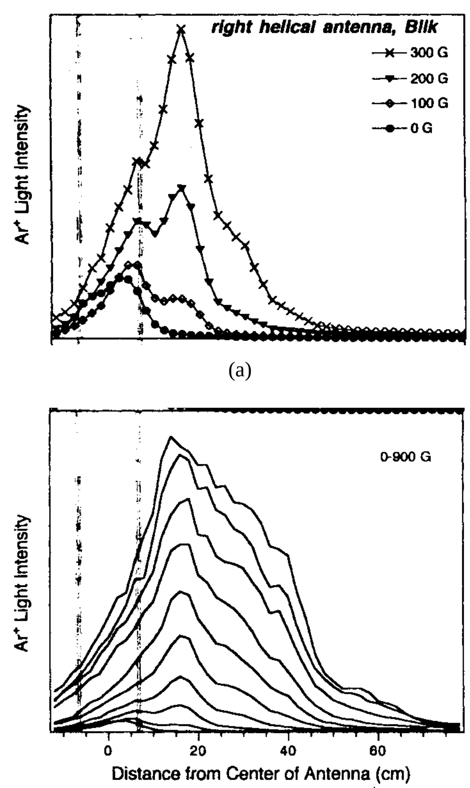
(a) (b)

Fig. 7. (a) Density jumps as B_0 is increased (from [5]) and (b) density jump as $P_{\rm rf}$ is increased [115].

caused by the higher collision frequency but may also have been an artifact of the floating-potential method used for RF compensation, as was later explained by Chen [155]. The effect of Landau damping on energy transfer has not been proved. Indeed, the most recent results [61] show that both the damping rate and the ionization rate can be explained on the basis of collisions alone. Thus, the validity of the Landau damping hypothesis is not yet resolved. There is ample evidence for a fast electron population, but it has yet to be shown that this is either caused by wave-particle acceleration or is relevant to the overall energy balance of a helicon discharge.

VI. Discharge Initiation

In Section II, it was mentioned that the helicon dispersion relation was obeyed in helicon discharges above a certain threshold. It is often observed, especially with plane-polarized antennas, that the plasma density takes one or two sharp jumps as either the RF power or the magnetic field is increased. Examples are given in Fig. 7. Ellingboe and Boswell [26] have studied this phenomenon extensively and have attributed the behavior to the coupling mechanism. Initially, at low powers or fields, the electrostatic potentials on the antenna are capacitively coupled to the gas near the walls and break it down into a weakly ionized plasma. As $P_{\rm rf}$ or B_0 is raised, the fluctuating magnetic field of the antenna induces large electric fields in the interior of the tube, and the density jumps upward. At this point, however, the density is still not high enough to satisfy the dispersion relation of (2), and what one has is a nonresonant inductively coupled plasma (ICP), which is commonly used in plasma processing. Finally, when (2) is satisfied, a helicon wave is generated, and the density takes another jump as the helicon ionization mechanism, whatever it is, takes over. The benefits of a helicon discharge over an ICP are dramatically demonstrated by measurements, shown in Fig. 8, of the ionized argon light as B_0 is raised from 0 to 900 G [61]. At B_0 = 0, a faint discharge is localized to the antenna region. Above about 300 G, a helicon peak rises up on only one side of the antenna—in the direction of propagation



(a)

(b)

Fig. 8. Axial distribution of ionized argon light intensity at various magnetic fields [61]. The bottom four curves in (b), are the same as those in (a); the curves are 100 G apart. The data are from a 5-cm diam discharge with 2 kW of 27.12 MHz power.

of the $m = +1$ mode. Eventually, this peak grows to 20–30 times the height of the ICP peak.

Boswell *et al.* have made time-resolved measurements in pulsed helicon discharges [115], [7], [33], giving further insight into the startup process. They found that a short burst of 200 eV electrons is generated initially, and these lead to rapid ionization of the gas. Acceleration of the electrons was attributed to a multipacting mechanism, presumably during the time when capacitive coupling was operative. Only after 70 ns did the discharge reach a steady state.

VII. Discharge Equilibrium

The study of helicon *discharges* is considerably more difficult than the study of helicon *waves*, since the processes of ionization and particle and energy transport have to be addressed. The problem can be simplified by dividing the plasma into the *near-field region* under and within a wavelength of the antenna, and the *downstream region* farther away. A survey of the parameters n_0, KT_e, and plasma potential V_s in the downstream region was made recently by Sudit *et al.* [61], [16], [17]. The temperature KT_e was found to peak $\frac{1}{2}$–1 wavelength past the antenna, and the discharge beyond this peak was defined as the downstream region. Since the wave amplitude was too small downstream to contribute to heating the electrons, the monotonic decay of KT_e downstream could be explained quantitatively by classical heat conduction balanced by electron energy loss from inelastic collisions with ions. The plasma density, surprisingly, rose away from the antenna to broad peak about 50 cm downstream before finally decaying (Fig. 6). The rise in density was consistent with the approximate constancy of $n_0(KT_e - e\phi)$ along the field lines and was therefore caused by the drop in KT_e.

With local measurements of KT_e and n_0, it was possible for Sudit and Chen [61] to calculate the radial losses due to classical diffusion at each position z, and compare those with the local ionization rate from the thermal electron distribution. In contrast with the B_0 = 0 ICP case, it was observed at high magnetic fields that the plasma was lost downstream from where it was produced, showing the effectiveness of magnetic confinement of the electrons. Agreement between calculated plasma production and loss integrated over z was achieved to within a factor of two. However, this degree of agreement was achieved only when a depletion of neutral density by a factor of 4–5 was taken into account. This phenomenon, first reported by Boswell [5], is caused by the gas-pumping effect [61], [33] in which the neutrals are ionized and then accelerated to the ends of the tube by presheath potentials of the order of $\frac{1}{2}KT_e$. Helicon discharges can be so dense that the neutral gas is completely removed from the center, and the plasma is 100% ionized there [5].

An interesting conclusion can be drawn from calculations of power balance. If one assumes classical electron heat conduction, the slope of $KT_e(z)$ at the boundary between the near-field and downstream regions gives the heat flux. It was found that this amounted to about 70% of the total RF input power, leaving only 30% to be transported backward to the short end of the tube and to be deposited by the waves in the downstream region [61]. This means that the RF power is mostly absorbed in the near-field region, whose length of about 20 cm is orders of magnitude shorter than the collisional or Landau damping length of helicon waves. Clearly, the near-field region holds the clue to the efficiency of helicon discharges. Ellingboe and Boswell [26] have made detailed measurements of the wave fields in this region. Exactly what happens here is not yet known definitively but will be discussed in Section IX.

VIII. Finite Mass Effects

When B_0 is sufficiently low, neglect of the quantity $\delta \equiv \omega/\omega_c$ eventually fails. Finite electron mass effects had been considered by early workers [143], [162], [152], [128], but it was only in recent work [166], [179], [129], [157], that we came to realize that the effects cannot always be neglected even if $\delta \ll 1$. There are actually several different effects. Electron inertia parallel to B_0 causes E_z to be nonzero even in the absence of collisions. This complicates the boundary conditions that have to be satisfied, but collisions are usually frequent enough to mask the effect of parallel inertia. Electron perpendicular inertia gives rise to another branch of the dispersion relation involving electron cyclotron waves. As we shall see, these waves are strongly coupled to helicon waves and provide much of the damping. At frequencies near ω_c, cyclotron damping provides another possible collisionless damping mechanism; this has been treated theoretically by Harvey and Lashmore–Davies [168] but studied experimentally in helicon plasmas only in an early paper by Christopoulos *et al.* [20].

When $\delta \neq 0$, (1) is replaced by

$$\delta\nabla \times \nabla \times \boldsymbol{B} - k\nabla \times \boldsymbol{B} + k_w^2\boldsymbol{B} = 0. \tag{6}$$

Since the differential equation is now of second order, a new wave appears: this is an electron-cyclotron wave, which, in a finite cylinder, is generally known as a Trivelpiece–Gould (TG) mode. We now have a helicon (H) wave, described by the last two terms of (6), coupled to a TG wave, described by the first two terms. In most previous work, helicon theory did not treat finite electron mass, and quasielectrostatic Trivelpiece–Gould theory omitted the electromagnetic effects incorporated in the last term, which is proportional to $\boldsymbol{\nabla} \times \boldsymbol{E}$. Factoring (6) yields [128] $\boldsymbol{B} = \boldsymbol{B}_1 + \boldsymbol{B}_2$, with $\boldsymbol{B}_1$ and $\boldsymbol{B}_2$ satisfying

$$\nabla \times \boldsymbol{B}_1 = \beta_1\boldsymbol{B}_1, \qquad \nabla \times \boldsymbol{B}_2 = \beta_2\boldsymbol{B}_2 \tag{7}$$

where β_1 and β_2 are the roots of

$$\delta\beta^2 - k\beta + k_w^2 = 0. \tag{8}$$

Thus, the β_1(H) and β_2(TG) roots are given explicitly by

$$\beta_{1,2} = \frac{k}{2\delta}\left[1 \mp \left(1 - \frac{4\delta k_w^2}{k^2}\right)^{1/2}\right]. \tag{9}$$

The k–β relationship of (8) is illustrated in Fig. 9. For given $k > k_{\min} = 2k_w\sqrt{\delta}$, the H branch has the smaller value of β, and the TG branch the larger value. Each branch has the wave profiles of (4), but with $T_j^2 = \beta_j^2 - k^2$. At large B_0, β_2 is much larger than β_1, and we see that the TG wave has short radial wavelength, while the H wave occupies the bulk of the plasma. These waves are coupled at the boundary, but since the TG mode is heavily damped, it is essentially a surface wave at large B_0, and the two waves are almost independent. However, as B_0 is decreased, the values of β_1 and β_2 approach each other, and the two branches become well

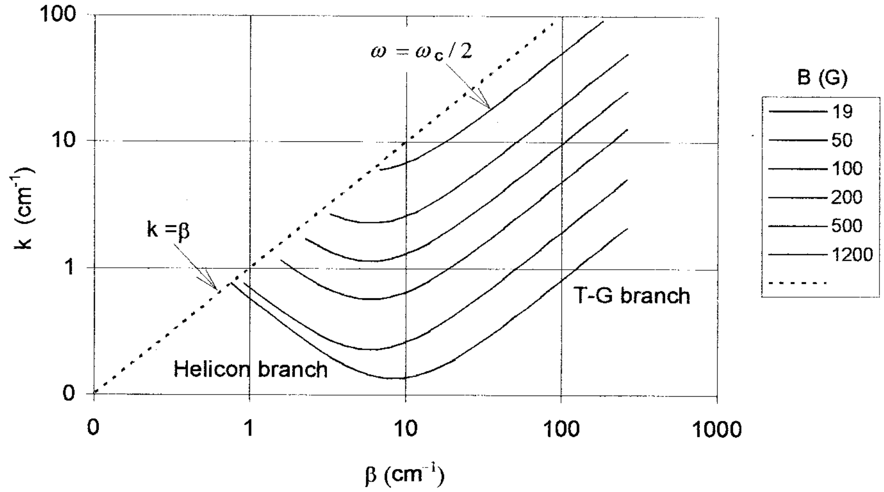


Fig. 9. Relation between parallel wavenumber k and total wavenumber β for a 2.5-cm radius discharge at 27.12 MHz and various magnetic fields. The diagonal line is the limit $\beta = k$ [129].

mixed throughout the plasma. When δ increases beyond 1/2, the H branch becomes evanescent. As B_0 approaches zero, both waves become evanescent coalesce into the skin-effect fields of the ICP, modified by the dc magnetic field.

These coupled modes have been seen by Cui and Boswell [21], who varied δ by increasing the frequency. A number of authors have seen density peaks occurring at low B-fields where the TG mode should be important [12], [23], [95]; however, why the density should increase there depends on the coupling efficiency, which will be discussed below.

With light gases at high magnetic fields or low frequencies, ion motions become important. The relevant small parameter here is $\gamma \equiv \Omega_c/\omega$, where Ω_c is the ion cyclotron frequency. A simple way to take finite-γ effects into account is to replace δ in (6) by $\varepsilon \equiv \delta - \gamma$, giving [128]

$$(\delta - \gamma)\nabla \times \nabla \times \boldsymbol{B} - k\nabla \times \boldsymbol{B} + k_w^2 \boldsymbol{B} = 0. \qquad (10)$$

We see that when $\gamma > \delta$, which easily happens with elements like hydrogen or helium, the "slow wave" changes from an electron wave to an ion wave. The transition occurs when $\omega = (\omega_c \Omega_c)^{1/2}$, at which point only the helicon branch exists. Peaks in density occurring at frequencies near the lower-hybrid frequency in light gases were first seen by Shoji [56] and recent measurements have been made by Sakawa *et al.* [51] and Yun *et al.* [66]. These peaks are not predicted by (10), which neglects dissipation. Explanation of these ion effects awaits computations which include damping and antenna coupling.

IX. Antenna Coupling

Several types of antennas have been used for applying RF fields to helicons. Until recently, the double-saddle coil antenna was used almost exclusively by Boswell and co-workers [4]–[8], etc. The plane-polarized Nagoya Type III antenna [85], [88] was adopted by Chen and co-workers [10]–[17], etc. Twisted Nagoya Type III antennas (helical antennas) of either RH or LH polarization were first used by Shoji [56]–[58] and later adopted by Chen *et al.* These antennas all have $m = \pm 1$ azimuthal symmetry, and the current return is via circular end-rings. The $m = 1$ end-rings alone have been used recently by Ellingboe and Boswell [26] and named a double half-turn antenna. Pure helices without end-rings have been used by Kim *et al.* [35], who have also tried $m = 2$ straight antennas [34]. In industry, symmetric $m = 0$ antennas have been developed and marketed successfully [94], [95], [60]. Single- and double-loop antennas have been investigated by Jiwari *et al.* [28]. Tests of the relative efficiency of different antenna designs have been made by a number of authors (e.g., [12], [26]).

Understanding of Nagoya Type III antennas came from work on ion-cyclotron isotope separation at TRW, Inc., in the 1970's. The strong magnetic field prevents electrons from short-circuiting the antenna field, and the skin depth is increased from c/ω_{pe} to c/Ω_{pi}, which is larger than the plasma radius. Furthermore, the $m = \pm 1$ geometry and finite $k_{\|}$ act together to convert the electromagnetic field of the antenna into an internal electrostatic field in the plasma [73], [12]. Also for this project McVey wrote the ANTENA code [83], [84], which is commonly used to model these antennas and their coupling to nonuniform plasmas.

The most complete experimental study of the near-field region where antenna fields couple to the plasma was made by Ellingboe and Boswell [26]. Using a double-half-turn antenna, they could clearly see sudden changes in the two-dimensional wave patterns as $P_{\rm rf}$ was raised and the coupling changed from capacitive to inductive, and finally to helicon waves. Using a double-saddle-coil antenna, they found that the fields under the antenna were standing waves, which became traveling waves as one moved from the antenna into the downstream region. This led them to propose that the coupling mechanism is the trapping of electrons in the large fields under the antenna, followed by acceleration of the electrons to ionizing energies as they moved into the traveling wave region. This model was further elaborated on by Degeling *et al.* [23], who showed that the measured phase velocity agreed with that for the most efficient ionization produced by electrons trapped from the thermal distribution, and furthermore that the scaling of density with $P_{\rm rf}$ agreed with the trapping model. Traveling waves can be produced even under the antenna if phased bifilar antennas are used [64], [45], and one might expect that trapping would be even more efficient in this case.

The Nagoya Type III coupling mechanism mentioned above involves only the antenna currents parallel to $\boldsymbol{B}_0$, and these are expected to excite wavelengths λ which are twice the antenna length L, and this $\lambda = 2L$ component is enhanced by the end-rings which provide the current return path. If the Type III antenna is helical, however, the end-rings can generate a $\lambda = L$ component. This was observed by Light *et al.* [42] and has been confirmed using two double-half-turn antennas by Ellingboe *et al.* [26].

On the theoretical side, a number of authors have made numerical calculations of antenna coupling; for instance, Fischer *et al.* [166], Mouzouris and Scharer [172], Molvik *et al.* [47], and Borg *et al.* [144], [145], [169], as well as the Boswell, Chen, and Shoji groups. A series of papers by Shamrai *et al.* [175]–[179] and concurrent work by Arnush and Chen [157] have recently thrown new light onto this problem as a result of the inclusion of finite electron mass effects. Since the TG mode is highly damped near the surface of the plasma, a new mechanism is available for efficient absorption of RF energy which is not available to other RF plasma generators. Because of the magnetic field's effect of electron motions, the antenna field can penetrate deeply into the plasma and couple to the weakly dampened helicon mode there. The H mode then mode-couples via the boundary condition to the TG mode, which is strongly damped as it propagates inward, depositing its energy into the outer layers of the plasma. Computations [138] show that the energy deposition has two peaks, one near the axis due to the H mode, and a larger one near the boundary due to the TG mode. Indeed, many authors have reported seeing, at low magnetic fields, a ring of high density at large radii (e.g., [23], [51]). In recent computations, it has been found that the antenna couples much more strongly to the $m = +1$ mode than to the $m = -1$ mode, perhaps because the $m = -1$ mode, with a narrower wave profile, has less amplitude at the edge for coupling to the dissipative TG mode there. This would explain why Shoji, in 1992, observed that an $m = +1$ helical antenna produced a dense, narrow plasma near the axis, while an $m = -1$ antenna filled the chamber with a lower density plasma, in spite of the fact that it has the narrower wave profile. It was simply that the coupling to $m = -1$ was so poor that only nonresonant ionization was produced. The validity of the mode-coupling absorption mechanism will no doubt be tested definitively in the next few years.

X. Industrial Applications

During the past decade, a large number of papers have appeared regarding the use of helicon sources in actual plasma processing applications [89]–[123], and the number is increasing rapidly. Review of this work is beyond the scope of this paper; however, some generalizations can be made. As in all high-density plasmas, the neutral pressure is so low that very few particulates are generated, and the problem with contamination that plagues capacitive discharges is not present in helicon plasmas [117]. A second problem, that of charge damage to thin oxide layers, is alleviated by the good uniformity and low magnetic field at the wafer surface that can be achieved with helicon sources. In silicon etching, the plasma plays a dual role and therefore has to be carefully formulated. As shown in early experiments on pulsed plasmas by Boswell and Porteus [92], it is neutral Cl (or F) that does the actual etching. For this purpose, the plasma needs only break up the molecules of the fill gas into atomic form. Ions are not wanted here, and in fact a low electron temperature is desired so that there is as little ionization as possible after the Cl atoms are formed. On the other hand, trenches and holes with straight walls can be etched only if anisotropy is provided by unidirectional ions accelerated in a planar sheath. These prepare the surface at the bottom of a trench so that Cl atoms can attack it efficiently in a well-known symbiotic process. For this purpose, charged particles are obviously needed. Thus, an ideal plasma for etching would have low T_e, many neutral atoms, and just enough ion flux to give directionality. A higher density plasma would increase the etch rate, but too high a density would overheat the wafer, increasing the cooling requirements, or make the exposure time uncontrollably short. There are many other procedures besides Si etching: oxide etch, metal etch, metal or oxide deposition, photoresist stripping, etc. Each of these requires a different type of plasma. Thus, a good plasma source for industrial use should have the flexibility to operate under a wide range of conditions.

The further development of helicon sources is proceeding on several fronts. First, the use of magnetic fields below 100 G is emphasized because of their low cost and the fact that densities in the 10^{13} cm^{-3} range, achievable at high fields, are not yet required. Second, the move toward silicon wafers larger than 30 cm in diameter, as well as flat-panel display substrates larger than 50 cm on a side, has triggered experiments with larger diameter and multiple-tube helicon sources. Finally, etching with negative ions is expected to decrease oxide damage and notching effects; and, consequently, there is experimentation with pulsed discharges, in the afterglow of which KT_e can decay to values consistent with electron attachment to form Cl^- or O^- ions [109].

The worldwide interest in helicons during the past decade has resulted in a much better understanding of the physical mechanisms than was available in the early years. During the next decade, we expect that these ideas will be converted into practical devices for technologies of the future.

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Experiment

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