Resonance absorption in a microwave plasma interaction

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A detailed study of resonance absorption and the resultant plasma wave electric field evolution is reported. At sufficiently low input power levels, convective saturation of the localized electric field is observed. At higher input power levels, wave-breaking saturation begins, resulting in hot electron production. The transition from the convective to the wave breaking regime occurred at much lower powers than expected. The temporal and spatial evolution of the harmonic fields of the excited plasma wave as well as the quasisteady magnetic field were also investigated and compared with existing theories.

I. INTRODUCTION

Resonance absorption in an unmagnetized, inhomogeneous plasma occurs when a plane electromagnetic (EM) wave is obliquely incident with the wave electric field polarization in the plane of incidence (P polarization). This wave will be cut off at the position given by \( \omega_c(z) = \omega_b \cos \Theta \), where \( \omega_c(z) \) is the local plasma frequency, \( \omega_b \) is the incident wave frequency, and \( \Theta \) is the angle of incidence. However, the evanescent field of the EM wave can penetrate to the higher density region and resonantly excite an electrostatic (ES) wave at the critical density layer (linear electromagnetic wave conversion to an electron plasma wave), thereby transferring energy from the EM wave to the plasma. This linear mode conversion has been studied extensively in the past, both analytically, 1-3 numerically, 4 and experimentally in microwave-plasma interactions.5,7

For the case of a cold plasma with a linear density profile \( n_e(z) = n_0(1 + z/L) \), the solution for the electromagnetic wave is given by an Airy function in the underdense region, while at the critical layer (\( z = 0 \)), the electric field becomes infinite in the absence of collisions. 1 Several phenomena can limit the amplitude of the electric fields driven by linear mode conversion. For nonzero electron temperature, the plasma waves excited near the critical layer have finite group velocities and convet energy out of the resonance region. The eventual balance between incident electromagnetic waves depositing energy in the Langmuir waves, and loss due to convection can limit the field amplitude. Other field-limiting mechanisms include particle collisions and wave breaking. The enhanced fields can be much larger than the incident and can lead to density profile modification, magnetic field generation and production of suprathermal electrons. These effects in turn may modify the resonance absorption and field amplitude limiting processes. For the laser fusion application, resonance absorption can be important; particularly for the case of higher laser irradiances and longer laser wavelengths where inverse bremsstrahlung is reduced. There is thus interest in it as an absorption mechanism and as a potential source of suprathermal electrons with accompanying target preheat. 2

Herein, we present experimental results where pulsed microwaves are obliquely incident on an inhomogeneous, near critical density plasma. The plasma target is effectively collisionless so that convection and/or wave breaking would be expected to limit the linearly converted wave amplitude. In particular, one would expect convection to dominate at lower power levels while wave breaking should dominate at higher powers. Parametric studies in the experiments (which included measurements of the enhanced fields and associated hot electrons) indeed show the transition between these two amplitude limiting phenomena. However, the transition from convection to wave breaking occurred at a much lower power than one would predict from theoretical models utilizing fixed ions. These results imply wavebreaking has a predominant role over a broader parameter range than expected.

The organization of this paper is as follows. Section II contains a brief review of the theoretical predictions on the saturation of resonance absorption by cold plasma wavebreaking. A description of the experimental apparatus and measurement techniques employed in these studies is given in Sec. III. The actual results are contained in Sec. IV together with comparisons with theory. Finally, Sec. V contains a summary of the principal conclusions of this work.

II. THEORY

In this section we summarize some theoretical results on saturation of the linearly converted wave amplitude. The simplest case is to account for convection of wave energy due to finite electron temperature. The plasma waves excited near the critical layer transport energy out of the resonance region with a velocity comparable to the plasma wave group velocity giving rise to a finite-steady-state field amplitude. In this case, the amplification factor for a linear density profile is given by

\[
E_s(z = 0)/E_0 = \phi(\tau)/(2\pi\rho)^{1/2}(\beta/\rho)^{3/2},
\]

(1)
where $\phi(\tau)$ is the resonance function,$^1$ $\tau = (k_0 L)^{1/3} \sin \Theta$, $\rho = k_0 L$, $\beta = u_e/c$, $u_e$ is the electron thermal velocity, $k_0$ is the vacuum wavenumber of the incident wave, and $E_0$ is the vacuum field intensity. The characteristic time for reaching this steady state is $t_c = (L / \lambda_0)^{3/2}/\omega_0$.

The amplitude of a strongly driven plasma wave can become comparable to the kinetic energy of the electrons moving at the wave phase velocity even if the field intensity $\eta_0 = (E_0^2/8\pi n_k T_e)$ is weak ($<1$). The electrons can then be trapped in the wave potential well with energy transfer occurring as the electrons bounce back and forth in the potential well and are accelerated. In the case of a strongly localized large amplitude oscillation such as occurs in the resonantly driven case or solitons, electrons can be accelerated through the wave in one oscillation period. This process is called wave breaking.

Let us consider the case of cold plasma wave breaking associated with the resonant absorption of a monochromatic pump (frequency $\omega_0$) of sufficiently low intensity that density profile modifications can be neglected. The electron displacement is governed by$^4$:

$$\delta + \omega_0^2 z_0 \delta = \delta_p \delta_0 \omega_0 e^{i\omega t},$$

where $\delta = eE_d/m_0 \omega_0^2 \sim L$, $\omega_0^2(z_0) = \omega_0^2(1 + z/L)$, and $z = z_0 + \delta(z_0)$. The driving electric field $E_d \delta$ is related to the vacuum field $E_0$ by $E_d = E_0 \delta(z)/(2\pi k_0 L)^{1/2}$. Wave breaking occurs when $\delta \delta / \partial z = -1$ which yields the wave-breaking time

$$t_{\delta} = (1/\omega_0)(8L / \delta_0)^{1/2} \sim P_o^{-0.35},$$

where $P_o$ is the incident rf power. Electrons are ejected with velocity $(2\omega_0^2 L \delta_0)^{1/2}$, which results in maximum energy $\epsilon_m = (1/2)m_0 \delta^2$ given by

$$\epsilon_m = m_0 \omega_0^2 L \delta_0 \sim P_o^{0.3}.$$

Equation (4) also implies that the wavebreaking electric field intensity $E_d = \epsilon_0^{1/2} \sim P_o^{0.5}$, where

$$E_d = (m_0 \omega_0 \epsilon_0)^{1/2} \sim (2L m_0 \omega_0^2 e E_0)^{1/2}.$$

Computer simulations have indicated that the suprathermal electron temperature $T_{he}$ is proportional to $\epsilon_m$, although the thermalization mechanism is still under discussion.$^{11}$

From this simplified analysis, one might expect convection to dominate when the wavebreaking time is longer than the convection time. With fixed ions, the convection time is only dependent on the electron temperature and the density scale length. For conditions where the electron temperature varies slowly, or is independent of the driver intensity, the convection time is independent of intensity. On the other hand, the wave-breaking time is a monotonically decreasing function of driver amplitude. Thus one would expect wave breaking to dominate at the higher intensities. However, as we shall see later, experimentally we observe that the transition to wave-breaking saturation occurs at a significantly lower power level than this simple physical picture predicts.

To understand this discrepancy we performed extensive experimental studies of the microwave-plasma interactions where the incident power level was varied over a broad range. This allowed us to access regimes extending from where the enhanced fields would produce negligible effects to regimes where the localized pressures associated with the enhanced fields exceed that of the background plasma and significant density perturbations would be expected. In addition, we performed a number of one-dimensional ESII computer simulations of the interaction using a capacitor plate model with moving ions.$^{12}$ These qualitatively reproduced many of the features observed experimentally.

III. EXPERIMENTAL ARRANGEMENT

The experiments were performed in a cylindrical, unmagnetized plasma (60 cm diameter, 80 cm length) produced by a pulsed argon filament discharge. Typical plasma parameters during the filament discharge are: maximum electron density $n_e = 1.7 \times 10^{11}$ cm$^{-3}$, electron temperature $T_e \sim 2.5$ eV, and ion temperature $T_i \sim T_e/10$ for an argon neutral gas pressure of $4 \times 10^{-4}$ Torr. In order to avoid the complications associated with the presence of the primary ionizing electrons, most of the experiments reported herein were performed in the plasma afterglow following the plasma discharge pulse. The plasma parameters are similar to those given above except that the electron distribution function is completely Maxwellian with $T_e \sim 1$ eV. The plasma density profile is approximately linear along the chamber axis with $L / \lambda_0 \sim 5 – 10$, where $\lambda_0$ is the vacuum wavelength of the incident electromagnetic wave. In addition to this axial density gradient, there is a weaker radial density gradient with $L_r / \lambda_0 \sim 20$. Pulsed microwave radiation ($\omega_0/2\pi \sim 3$ GHz) of typical rise time 10–20 nsec is launched along the chamber axis using a mesh horn. Pulse rise times as short as 1 nsec ($\sim 3$ rf periods) were available when required.

The angle of incidence $\Theta$ (measured at a radial position of approximately 3 cm from the axis) of the incident microwave radiation is $\Theta \sim 10^\circ – 14^\circ$ with respect to the central plasma density gradient. This was determined in several different ways as discussed below.

Electric fields were measured by small movable monopole (single wire) antennas. The signals were transported via 50 ohm coaxial cables to calibrate receivers which discriminated between signals near the fundamental microwave frequency and the harmonics. No significant difference was observed in the signals when using dipole rather than monopole probes or when the antenna wire bias was varied. Although the probes were small, one must, of course, consider the possibility that their presence may modify the interaction when inserted near the resonant region. Correlated remote measurements were thus made of plasma electrons accelerated by the electric fields using retarding field energy analyzers. These energy analyzers were located well away ($\sim 10$ cm) from the enhanced field regions. The presence (or absence) of the antenna probes had no appreciable effect on the accelerated electrons observed by the energy analyzers.

The magnetic field generated by the incident microwave radiation is measured by magnetic loop coils (ten turns, 0.5 cm and 1.0 cm diameter) covered with thin copper foil to reduce the direct induction from the intense microwave source. The signal output is amplified by a fast rise time

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differential amplifier and is passively integrated or digitized then integrated numerically.

IV. EXPERIMENTAL RESULTS
A. Plasma parameters

As mentioned earlier, the work reported herein is concerned primarily with the interaction of low intensity \((\eta_0 \sim 3.5 \times 10^{-3})\) microwave radiation with an inhomogeneous plasma. Particular attention has been paid to the first 0.5 \(\mu\)sec of the interaction \((\omega \tau \sim 10^8)\).

Resonance absorption is dependent upon the angle of incidence of the electromagnetic wave with respect to the density gradient. In our experiment, the microwaves were launched by a horn along the central \((z)\) axis of the cylindrical chamber. The angle of incidence inferred from the measured axial and radial density profiles is approximately 14°. The geometry is complicated by the fact that the interaction occurs in the quasinear field of the horn antenna with the radiation diverging at an angle comparable to the angle of incidence. Nevertheless, the behavior of the electromagnetic wave–plasma interaction near the chamber (and horn) axis appears similar to that of a planar electromagnetic wave.

Figure 1 shows a spatial scan of the electric field at a low-power level. The angle of incidence can also be unfolded from the separation between the two rightmost peaks (the farthest one corresponding to the critical layer) since the electric field intensity distribution is describable by an Airy function as a result of the linear density variation. Utilizing the properties of the Airy function yields the result that

\[ \Theta = \sin^{-1}\left(\frac{D}{L}\frac{d}{L}\right)^{1/2}, \]  

where \(D\) is the separation between the two peaks and \(d = -\left(L/k_n\right)^{1/2}\). Using the data of Fig. 1 together with the measured scale length \(L \sim 50\) cm yields \(\Theta \approx 8°\) in reasonable agreement with the angle of incidence determined by the density profile measurements. An additional confirmation is obtained from absolute measurements of the growth of the spontaneous magnetic field associated with resonance absorption. Specifically, the initial growth rate near the critical layer is given by

\[ \dot{B}_y \sim 4\pi n_e m_e \eta_0 (nk_T) \sin \Theta \cos \Theta, \]  

where \(A \sim 0.5 \) is the angular-dependent absorption coefficient. By measuring \(B_y\), \(n_e\), \(T_e\), etc. we can calculate \(\Theta\) and typically obtain \(\Theta \sim 10°-14°\).

At the higher power levels, we observe the growth of localized large amplitude electric fields, \(E^2\), near the critical layer together with the production of hot electrons, \(I_{hot}\), and quasistatic magnetic fields, \(B_y\) and \(B_x\). Figure 2 shows typical oscilloscope traces for an incident rf power level of 500 W.

Density profile modifications can affect resonance absorption and thus it is necessary to monitor the plasma density profile in order to assess their importance. Figure 3 shows typical density perturbation data obtained by sam-

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**FIG. 1.** Typical spatial scan of the \(z\) component of the electric field intensity at low input power (\(P_e = 10\ W\)) and sampled 390 \(n\)sec after turn-on of the rf field.

**FIG. 2.** Time histories of rf pulse (\(P_e = 500\ W\)), electric field intensity \(E^2\), hot electron current (\(\sim 20\) ev), \(I_{hot}\), time derivative of quasisteady magnetic field \(B_y\) and quasisteady magnetic field strength \(B_x\).
plunging a Langmuir probe signal approximately 100 nsec after
turnoff of the rf pulse whose duration for each incident pow-
ner level was slightly longer than the measured wave-breaking
time. By making the measurements in the absence of the rf,
spurious probe effects such as rectification of the rf are elimi-
nated. Notice that even at 500 W ($\eta_0 = 3.5 \times 10^{-3}$) the
modifications are rather modest ($\delta n/n_0 \sim 12\%$). The role
such density profile modification plays will be discussed in a
later section.

B. Observation of wave convection saturation

As can be seen in Fig. 2, the electric field initially begins
to grow, then exhibits one or two inflection points or shoul-
ders before rising to its peak value and subsequently decreas-
ing to a significantly smaller value. As we shall see in the
following, we have identified convective saturation followed
by cold plasma wave breaking, respectively, as being responsi-
ble for the above temporal behavior of the electric field.

First of all we begin with a study of the initial saturation
(shoulder) of the electric field growth which we have identi-
fied as convective saturation. In Fig. 4(a) we plot the magni-
tude of the electric field intensity (measured at the first
shoulder) as a function of the incident rf power. The depen-
dence of electric field intensity on input power is given by
$E_0^2 \propto P_0^{0.81}$ close agreement with the anticipated linear
scaling with incident power for convective saturation [see Eq.
(1)]. The convective saturation time should be independent
of incident power. As shown in Fig. 4(b) the measured satu-
ration time is in fact essentially independent of rf power for
$P_0 \sim 4$ W. The experimental saturation times of $\sim 180$ nsec
are in reasonable agreement with the calculated convective
saturation time $t_e$ for our experimental parameters of $\sim 75$–
100 nsec.

Referring back to Fig. 2, following the initial convective
saturation there is a further growth in the electrostatic field
intensity which saturates at later time with a significantly
larger amplitude. As discussed previously, one might expect
that the convective saturation mechanism would determine
the field amplitude since our pump intensities are quite weak
($3.5 \times 10^{-3} < \eta_0 < 3.5 \times 10^{-3}$), and appear to be far below
those given by the wavebreaking limit. However, if we calcu-
late the quiver velocity, $v_s$, caused by the associated convec-
tive saturated electrostatic field we find that $v_s > v_e$ even for
power levels as low as 50 W. For such velocities the convec-
tive saturation model is not expected to rigorously hold and
other saturation mechanisms such as particle trapping and
cold plasma wavebreaking might be expected to occur. In the
following, we discuss the supporting evidence for the
wavebreaking interpretation for the final saturation. How-
ever, before considering the region above 50 W, let us consid-
er the low power region ($P < 50$ W). Computer studies by
Morales and Lee have shown that where $\rho = (\omega L/v_e)^2
(E_0^2/2\pi n_e T_e) > 1$, density cavities and local trapping of
the high-frequency electric field can occur. For our plasma
conditions, when $\rho = 1$, the input power is 61 mW at the criti-
cal layer, and one expects to observe density profile distor-
tion. Therefore, pure wave convection is only expected to occur at
input power levels of less than 61 mW. A transition region
from pure wave convection to wave breaking should be ex-
pected for higher power levels.

In order to study the expected transition from pure conve-
tection saturation to wave breaking saturation, we per-
formed detailed studies in the region 10 mW < $P_0$ < 100 W.
The saturated electric field intensity is plotted as a function of
the rf input power level in Fig. 5. We can see that for
sufficiently low power ($P_0 < 2$ W), we do in fact observe a
scaling consistent with "pure" convective saturation and that
the saturation field intensity is linearly proportional to
$P_0$. We further observe that for $P_0 \approx 5$ W the saturated field
departs from the predicted convective saturation value and

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig3.png}
\caption{Density profile modifications at various incident power levels. Top trace is obtained without microwave power.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig4.png}
\caption{Electric field intensity and convective saturation time as functions of incident rf power: (a) initially saturated (inflection point) electric field intensity; (b) convective saturation time.}
\end{figure}
C. Observation of wavebreaking

For $P_o > 50$ W, the electric field amplitude appears to initially be saturated by wave convection, only to grow again at a later time. The continued growth implies that some mechanism impedes plasma wave convection and enhances the electric field strength. As predicted by Chen and Liu,\(^{15}\) when $\omega_p/\nu_e > 1$ and

$$\frac{E_d^2}{8\pi n_k T_e} > \frac{1}{8} \left( \frac{\Lambda_D}{L} \right)^2,$$

(7)

density profile modification will occur before convective saturation. The profile modification will slow down the wave convection and hence the convective saturation will occur at a delayed time. The continuous density profile modification will result in the localization or trapping of the electric field and the field will return to the growth phase. In fact, before convective saturation is reached we already have $v_w > v_e$ when $\eta_0 > 7 \times 10^{-4}$. This was verified both from hot electron measurements and electric field measurements. When the wave amplitude becomes sufficiently large, the wave breaks and hence the field intensity collapses. The measured peak amplitude (just before its collapse) appears consistent with the amplitude being limited by wave breaking. Figure 7 shows the measured time $t_b$ at which the maximum field intensity occurs as a function of the incident power. We see that $t_b \propto P_o^{-0.2}$ is in close agreement with the expected $P_o^{-0.35}$ scaling for cold plasma wave breaking [Eq. (3)]. In this region the measured profile modifications are modest ($\leq 14\%$) so that this agreement might be anticipated. Figure 5 again shows the scaling of the wave breaking electric field intensity with incident power above 60 W. The $P_o^{0.35}$ scaling for power levels below 600 W is again in good agreement with cold plasma wave breaking predictions.

The cold plasma wave breaking theory with which we compare the experimental results assumes that the background electron distribution is a $\delta$ function ($T_e = 0$). However, in a real experiment, there is a finite electron temperature. The effect of electron temperature on the wave-breaking amplitude was investigated theoretically by Kruer.\(^{16}\) He found that thermal effects increase the spatial extent of the field at wave breaking while decreasing its am-

![Graph showing the electric field intensity as a function of incident rf power.](image)

**FIG. 5.** Saturated electric field intensity as a function of incident rf power. "A" corresponds to the wave convection case and "B" to the wavebreaking case. Solid lines are the theoretical results with fixed density gradient and dashed lines are those with the measured local density profile modification included.

increases as $P_o^{1.2-1.8}$ until $P_o \approx 50$ W($v_w \approx v_e$). The initial saturation time (determined by the inflection point in the electric field growth) gradually decreases as seen in Fig. 4(b) in this power level. Above this power level, wavebreaking saturation appears as will be shown in the next section.

There are several phenomena which may be occurring in the transition region. First, the modest density perturbations can trap electron plasma waves thereby slowing down the loss rate and permitting a larger electric field build-up. The large amplitude plasma wave can trap more electrons which results in saturation of the electric field. Particle trapping effects may shorten the initial saturation time as seen in Fig. 4(b) for power levels $\approx 5$ W. Second, as mentioned earlier, the simple convective saturation mechanism becomes somewhat suspect for $v_w \approx v_e$. It is interesting that this transition occurs at a considerably lower power (50 W) than that (20 kW) predicted by the usual method of equating $E_e$ to $E^*$ at the transition. To further demonstrate the difference between the lower power behavior ($< 10$ W) and that observed at higher powers ($P_o > 50$ W), we compare the time evolution of $E^2$ in Fig. 6 ($P_o = 7$ W) with that shown in Fig. 2 ($P_o = 500$ W). Note in Fig. 6 the simple growth and subsequent saturation on the convective time scale with essentially constant amplitude over the duration of the rf pulse. The pronounced wave breaking feature of Fig. 2 is absent.

![Graph showing the time history of the electric field intensity near the critical layer at low power.](image)

**FIG. 6.** Time history of the electric field intensity near the critical layer at low power ($P_o = 7$ W).

![Graph showing the scaling of wavebreaking time as a function of incident power.](image)

**FIG. 7.** Scaling of wavebreaking time as a function of incident power.
amplitude. However, the plasma temperature only has a small effect on the wave-breaking amplitude (less than 40%) over a wide range of initial temperatures. Krueer also found that main body temperature does not strongly affect the suprathermal electron energy. He found that the enhanced fields were distributed over a broader area with increased main body electron temperature which compensated for the lower field amplitude in accelerating suprathermal electrons. Capacitor model simulations show that the maximum heated electron energy varied by less than about 10% when the background electron temperature was changed by a factor of 100.

Although the density profile is not strongly modified for power levels below 500 W, there is some local modification in the vicinity of the critical layer as shown in Fig. 3. Since the local density gradient scale length changes from \( \approx 50 \) cm to \( \approx 10 \) cm, it is appropriate to ask how accurately fixed profile wave breaking and convection theories describe the experimental conditions. The theoretical dependence on density gradient is weak with the wavebreaking and convective field amplitude scaling as \( E_x \propto L^{0.23} \) and \( E_z \propto L^{0.16} \), respectively. We used our experimentally measured local density gradient in the theoretical curves (dashed lines in Fig. 5) and the data were fitted at one point \( (P_0 = 213 \text{ W}) \) on the wave breaking theoretical saturation curve. As can be seen, the fit is extremely good.

D. Observation of ion wave streamers

When the microwave input power level exceeds 600 W, we observe a sharp transition to significantly more complicated behavior than described in the proceeding sections. Accompanying the growth and collapse of the electric fields we observe density disturbances (streamers) which propagate up and down the density gradient and originate in the region of maximum field intensity. These are apparently produced by the enhanced fields where the radiation pressure \( E^2 \) approaches and exceeds the background plasma pressure. The properties of the streamers were measured with Langmuir probes located remotely (\( \sim 2 \) cm) from the enhanced fields. Figure 8 shows the amplitude of the streamer as a function of rf pulse length. The streamer amplitude monotonically increases with pulse duration before reaching a saturated level. Interestingly, the pulse length where the streamer amplitude saturates roughly coincides with the wave-breaking time of the electric field. The signal from an antenna probe located in the center of the enhanced field region is shown for comparison in Fig. 8. The remote measurements of the streamers thus add confirmation to the wave breaking time of the enhanced fields measured with the antenna probes.

Not only do both the electric field intensity and wave-breaking time scale properly with incident power as predicted, but the measured wave-breaking times are in quantitative agreement with theory.\(^7\) Using our measured scalelength and angle of incidence, we estimate that \( E_x = 2.77 \text{ V/cm} \) at an incident power level of 500 W. This results in a predicted wave-breaking time of \( \sim 290 \text{ nsec} \) which is in reasonably good agreement with the measured value of 400–450 nsec.

E. Observation of higher harmonics

For \( P_0 < 50 \text{ W} \), the electric fields appear to be initially saturated by plasma wave convection as mentioned in Sec. 4.3. However, during the subsequent phase of the wave–plasma interaction, the density profile modification slows the plasma wave propagation and the field amplitude grows again until the wave breaks. Spatially resolved measurements of the plasma waves were made interferometrically by mixing a fixed reference probe signal with that from a movable probe. We find that the wave characteristics are quite different for the lower power and higher power cases. Specifically, at higher power levels we observe shorter wavelength components indicative of wave steepening and breaking.

As the amplitude of the electron oscillations increase during the approach to wave breaking, strong harmonic content in their velocity and associated electric field spectra is anticipated.\(^8\) We therefore set out to look for such harmonic content. Since the fundamental electric field has finite extent [see Fig. 9(a)], harmonic generation will occur in a region of finite extent. The spatial distributions of the second and third harmonics were investigated. It was found that the spatial extent of the second harmonic is comparable to or even narrower than that of the fundamental as is expected, but that of the third harmonic is much narrower as shown in Fig. 9(b).

Figure 10 shows the time histories of the total electric field intensity together with that of the second and third harmonics. We see that the peaks coincide but that the growth of the harmonics is delayed as expected until the waves nonlinearly steepen. The second harmonic field intensity is \( \sim 16 \text{ dB down} \) from the total field intensity while the third harmonic field intensity is \( \sim 41 \text{ dB down} \) from the total field intensity at the time of initial wave convection saturation.

It is of interest to compare the experimental results with theoretical calculations.\(^9\) However, since this theory is only valid for an electron oscillation amplitude \( \delta < 1/2k \), we have restricted our comparison to an early stage in the interaction when the oscillation amplitude is small. In the phase where
the plasma oscillation is limited by plasma wave convection, the electron density perturbation can be calculated as follows. First, the equation of continuity gives

\[ n_1/n_0 = (k/\omega)v_1. \]  
(8)

The plasma wave dispersion relation is

\[ \omega^2/k^2 = (3L/2\pi)v_e^2, \]  
(9)

where it has been assumed that the wavelength of the plasma wave \( \lambda_p \ll L \) and \( z \) is the distance from the critical layer in the underdense region. Thus,

\[ \xi = \frac{n_1}{n_0} = \left( \frac{2\pi}{3L} \right)^{1/2} \frac{v_1}{v_e}. \]  
(10)

Using the experimentally determined location for the maximum convectively saturated field, \( z \approx 0.4 \text{ cm} \), then \( \xi = 0.17 \) is obtained, where \( L = 100 \text{ cm}, T_e = 1 \text{ eV}, \) and \( P_0 = 500 \text{ W} \) are employed. Combining the above results with the theoretical predictions, the results in second and third harmonics field intensities that are 19 and 34 dB down, respectively, from the total field intensity in good agreement with the experimental observations. Here, we should point out, that strictly speaking this comparison is not valid since the theory assumed a cold plasma whereas experimentally we are concerned with a warm plasma. However, since the actual warm plasma dispersion effect is quite small \( (\Delta \omega/\omega = (3/2)^{1/3}kT_e/\omega_0 \approx 0.063) \) it can be expected to yield the approximate rate of harmonic generation.

Figure 11(a) shows that the wave-breaking time of the second harmonic (i.e., the peak time of the second harmonic electric field intensity) scales with incident power as \( P_0^{-0.29} \). To within the experimental error, this is in substantial agreement with the measured wave breaking time of the fundamental field \( t_b \propto P_0^{-0.29} \). The experimental results are also compared with the theoretical calculations obtained at the wave-breaking stage. Figure 11(b) shows the comparison of the theoretical calculation (solid line) and the experimental data (dots). The theoretical curve was calculated with the density perturbation \( n_1/n_0 \propto E_e \), where \( E_e \) is the convection-
limited plasma oscillation amplitude. The experimental data correspond to the second harmonic field intensities at the time the fundamental electric field is initially saturated by wave convection (refer to Fig. 10, the first peak of the electric field trace). As we can see in Fig. 11(b), agreement between theory and experiment is reasonably good.

F. Observation of hot electrons

Coincident with the electric field growth, a burst of hot electrons is observed (see Fig. 2). For the data shown in Fig. 2, the analyzing voltage was adjusted so that electrons with \( v \sim 4.5v_n \) were detected. Figure 12 shows the detected electron current versus analyzing voltage for various rf input power levels. The initial afterglow plasma is well described by a Maxwellian distribution with temperature of about 1–2 eV. When rf is applied we observe an increase in the bulk temperature as well as the growth of a hot electron tail with a distribution which can be described by a Maxwellian with temperature \( T_h \) up to a maximum energy \( \epsilon_m \) beyond which the distribution rapidly decreases. The data shown in Fig. 12 represent an average of many plasma repetitions with a sampling window of \( \approx 100 \) nsec width centered at the wave-breaking time. The scaling of the tail temperature \( T_h \) and cutoff energy \( \epsilon_m \) are shown in Fig. 13. A least-squares fit to the data yields \( T_h \approx P_0^{0.7} \) and \( \epsilon_m \approx P_0^{0.37} \). The latter is in good agreement with the predicted scaling of \( P_0^{0.5} \) from the cold plasma wave breaking model. Although thermalization is not covered by the wave breaking picture, it is of interest to note that \( T_h \approx P_0^{0.7} \) is predicted theoretically\(^{20}\) and obtained in simulations.\(^{11}\)

G. Observation of magnetic fields

As seen in the previous section, a burst of hot electrons and associated current is observed in our microwave plasma interaction. Such currents induce quasisteady state magnetic fields as have been observed in laser-plasma experiments\(^{21-23}\) as well as in microwave plasma experiments.\(^{15,24-27}\) The present investigation has concentrated on the study of the generation of quasisteady magnetic fields at relatively low intensities (\( \eta_0 \sim 10^{-2} \)) and at early times where resonance absorption instead of parametric instability\(^{12,28}\) would be expected to dominate.

In Fig. 2, an example of the quasisteady magnetic field is shown together with the growth rate \( B_v \). Here the temporal derivative of \( B_v \) is nearly proportional to the growth rate since the magnetic field increases linearly with time. We see that the maximum of the magnetic field growth rate critical is almost coincident with the maximum electric field intensity as well as the peak hot electron emission. This is expected if the local electric field gradient is sharp enough compared with the wave convection width near the resonance layer.

In Fig. 14, typical waveforms are shown obtained with a 2 \( \mu \)sec duration rf pulse as the intensity is varied. The magnetic field increases almost exponentially with time until \( \sim 1 \mu \)sec after turn on of the rf pump and then gradually saturates for larger pulse durations. In this power range, the magnetic field increases nearly linearly with the pump intensity. After turn-off of the rf pump it decreases almost exponentially, showing no evidence of an anomalous decay mechanism. Here, to ensure that the coil signal was a result of the B-field and not rf pickup, the coils were rotated through 180°.
FIG. 14. Time development of quasi-steady magnetic field produced by a short duration (~2 μsec) rf input pulse for various power levels.

to check that the sign of the field components changes appropriately.

For this experimental geometry, one expects that the $B_y$ component of the magnetic field should dominate if resonance absorption is the principal field generation mechanism. This is indeed observed with $B_y > B_x$, and $B_x > B_z$. The results shown in both Figs. 2 and 14 are for the $B_y$ component.

The magnetic field profiles are also measured two-dimensionally, along the axis of the plasma column (in the $z$ direction) and in the radial direction at constant $z$ position. An example of the field distribution in the $z$ direction is shown in Fig. 15 taken at various intervals in time for an rf pulse width of 1 μsec. As seen in Fig. 15, at earlier times up to rf turn-off ($t = 1$ μsec), there is a well-defined current sheet in the vicinity of the resonance layer. However, after rf turn-off, the main current sheet disappears and magnetic field line reconnection occurs resulting in localized bubblelike magnetic field shapes as seen from the data corresponding to $t \sim 2.0$ μsec. The growth rate of the magnetic field was also measured and found to be essentially in quantitative agreement with predicted value by Speziale and Catto. The initial growth rate near the critical layer is given by Eq. (6). Using the measured parameters, we predict a growth rate $B_x \sim 2.2 \times 10^{-4}$ G/μsec at $P_0 = 50$ W, which is in close agreement with the experimentally measured value of $\sim 4 \times 10^{-4}$ G/μsec.

The scalings of the initial growth rate and the saturated magnetic field were also measured. For $P_0 < 50$ W ($\eta_0 = 3.5 \times 10^{-7}$) both scale linearly with pump power as expected for resonance absorption produced magnetic fields. However, for incident power levels above 50 W, they scale as $B_x \propto P_0^{0.56}$. As seen in Fig. 5, the electric field initially saturates by wave convection and then grows again with a fast growth rate until it saturates by wave breaking. The transition occurred at $P_0 \sim 50$ W. Therefore, the fast growth rate observed in the magnetic field generation at higher incident power is also expected. An estimation of the saturation time will provide us with an approximate value of the saturated magnetic field. Speziale and Catto roughly estimate the magnitude of the saturated magnetic field strength by assuming that it grows with a linear growth rate until the convection loss time $\tau_c$ is reached beyond which time it is assumed constant in magnitude. The wave convection saturated quasi-steady magnetic field amplitude at the critical layer is given by

$$B_c = (A / \omega_c) u_0 E_0 \sin \Theta \cos \Theta \left( \frac{\omega_c}{L} / \omega_0 \right)^{1/3} \propto P_0 \quad (11)$$

where $u_0 = eE_0 / m_0$ is the quiver velocity in the incident microwave field. Using our experimental parameters, we find that $B_c (z = 0) \sim 5.1 \times 10^{-2}$ G. Unfortunately, this result is four orders of magnitude smaller than our experimental results. However, if we take into account the experimentally observed electric field strength, we find $B_c \sim 2.6 \times 10^{-2}$ G which agrees with the experimentally measured magnetic field $B_c \sim 3 \times 10^{-2}$ G at $P_0 = 17$ W. Thus the magnetic fields measurements are in agreement with the “anomalously” large electric fields which cannot be accounted for by convective saturation. An input power of 17 W is in the transition regime between pure convection and pure wave breaking as found from the electric field measurements.

When the oscillation amplitude of the plasma wave grows large enough, the wave breaking mechanism limits the field growth. Again, Speziale and Catto found the saturated field amplitude at the critical layer as

$$B_s = B_0 A \left( \frac{u_0}{c} \right) \left( \frac{\omega_c}{\omega_0} \right)^{1/2} \times \left( \frac{\omega_k L}{A \cos \Theta} \right)^{1/4} \sin \Theta \cos \Theta \propto P_0^{0.75} \quad (12)$$

This gives a saturation value $B_s = 1.47 \times 10^{-5}$ G for our experimental parameters. For an input power of 350 W, $B_s$ is found to be $1.2 \times 10^{-3}$ G, which is far smaller than our experimental results ($B_s \sim 0.1$ G with the same input power level). Using a simple model of dc magnetic field generation caused by resonance absorption together with the assumption $\partial E_x / \partial x \gg k_x E_x$, we obtain
FIG. 16. Space-time diagram of the current sheet near the critical layer obtained with the same conditions as in Fig. 15.

\[ B_b \approx (e/\imath m \omega_0) (E_b E_x) \]  

where \( E_b \) is the wave-breaking-limited electrostatic field amplitude. Therefore, we found that at an input power level of 350 W, the magnetic field amplitude at the critical layer is \( \sim 0.52 \) G which is larger than \( B_b \sim 0.1 \) G, perhaps indicating that other processes may be contributing to the field strength.

If we carefully look at the results shown in Fig. 15, it turns out that the current sheet moves down the density gradient. The zero crossing point of the magnetic field is plotted in Fig. 16, which shows that up to \( \sim 1 \) \( \mu \)sec after rf turnon the velocity is almost constant at \( 5 \times 10^6 \) cm/sec, but it gradually decreases to zero after rf turnoff (\( r = 1 \) \( \mu \)sec). The velocity in the early stage is considerably higher than the expected ion acoustic velocity of \( 2.0 \times 10^5 \) cm/sec obtained from the propagation of ion density streamers such as shown in Fig. 8. This might come from the expansion of the electron current sheet by the \( J_e \times B_z \) force generated by the self-field, which gives \( v_e \sim 1 \times 10^5 \) cm/sec. However, at present we do not possess a clear understanding of this mechanism.

V. SUMMARY AND CONCLUSIONS

In this paper, a detailed study of resonance absorption and the resultant plasma wave field evolution and associated phenomena were reported. The saturation mechanism(s) for these excited plasma waves were also studied in detail. The results of these experimental observations are summarized in Table I.

At sufficiently low input power levels, convective saturation of the localized electric fields produced by resonance absorption is observed. However, at anomalously low pump intensities, the electric field intensity departs from the expected linear scaling with incident power. Specifically, the saturated electric field begins to increase more rapidly with increasing input power than predicted by simple models. Magnetic fields generated through the microwave plasma interaction also showed a corresponding change in scaling at anomalously low power levels. This behavior can only be accounted for by more sophisticated models which account for localized density profile modification by the enhanced fields. The experimental results indicate that such complications cause phenomena associated with wave breaking to dominate at much lower power levels than expected. Detailed measurements in this low power wave breaking regime revealed that in addition to the fundamental electric fields, the quasistatic magnetic fields, density perturbations, hot

<table>
<thead>
<tr>
<th>( P_0 ) (W)</th>
<th>( \eta_0 ) (10^-5)</th>
<th>Electric Field Saturation Mechanism</th>
<th>Associated Phenomena</th>
</tr>
</thead>
<tbody>
<tr>
<td>5</td>
<td>3.5</td>
<td>Wave Convection</td>
<td></td>
</tr>
<tr>
<td>10</td>
<td>7</td>
<td>Transition Region</td>
<td></td>
</tr>
<tr>
<td>50</td>
<td>35</td>
<td>Wave breaking</td>
<td>hot electron production</td>
</tr>
<tr>
<td>100</td>
<td>70</td>
<td></td>
<td>quasi-statically magnetic field</td>
</tr>
<tr>
<td>500</td>
<td>350</td>
<td>Not Completely Identified</td>
<td>higher harmonics</td>
</tr>
</tbody>
</table>

TABLE 1. Summary of experimental observations.
electron generation, and harmonic field strengths had temporal evolutions and power scalings consistent with that expected from wavebreaking. Thus, despite the complication involved, in the physical mechanism, the results were surprisingly simple and consistent.

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