Acceleration and scattering of injected electrons in plasma beat wave accelerator experiments

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The results from experiments in which a two-frequency CO$_2$ laser is used to beat-excite large-amplitude, relativistic electron plasma waves in a tunnel-ionized plasma are reported. The plasma wave is diagnosed by injecting a beam of 2 MeV electrons and observing the energy gain and loss of these electrons, as well as the scattering and deflection of the transmitted electrons near 2 MeV. Accelerated electrons up to 30 MeV have been observed. The lifetime of the accelerating structure as seen by small-angle Thomson scattering is about 100 ps, whereas the injected electrons are seen to be scattered or deflected by the plasma for several ns, with diffuse scattering occurring 0.5–1 ns after forming the plasma wave and whole beam deflection occurring at later times. A simple model, which includes laser focusing, ionization, transit time, and relativistic saturation effects, suggests that the wave coherence may be short lived while the wave fields themselves persist for a longer time. This may be the reason for the disparate time scales between the Thomson scattering and the electron scattering diagnostic. The whole beam deflection may be evidence for a Weibel-like instability at later times.

I. INTRODUCTION

Plasmas may have useful applications in future accelerator technology. A relativistic electron plasma wave (where the Lorentz factor $\gamma_{\text{ph}}$ associated with the phase velocity of the wave is $> 1$) of amplitude $\epsilon$ propagating in a density of $n_e$ has a longitudinal electric field in GeV/m units given by $E_{\text{accel}} \approx 10^{-7} \epsilon \sqrt{n_e}$. For example, an $\epsilon = 30\%$ wave in a $10^{17}$ cm$^{-3}$ plasma would have a field of about 9.5 GeV/m. Thus, an electron could be brought up to an energy of, say, 1 GeV in a little over 10 cm. This rate of acceleration is about $10-100 \times$ higher than the state of the art in conventional radio-frequency (RF) structures, which is limited by RF breakdown. For the 1 GeV application, the 10 cm plasma and the “tabletop” (100 J, few ps) laser driver would replace 10–100 m of linac structure and the associated RF power.

In this paper, we discuss a series of experiments designed to demonstrate the feasibility of laser beat-wave-driven plasma waves as accelerating structures for electrons. The experiment has been discussed in great detail in previous papers, and so Secs. II and III will have only a brief description of the apparatus and some of the results, respectively. In Sec. IV we will introduce the simple model developed to gain some insight into the results. We discuss the experimental results in terms of the model in Sec. V and summarize in the final section.

II. EXPERIMENTAL SETUP

The experiment is shown schematically in Fig. 1. A two-frequency laser beam (10.591 and 10.275 pm in a two- or three-to-one ratio, typically) and a 2 MeV electron beam are both focused to the same point in a vacuum chamber filled with about 155 mT of hydrogen gas. The $1/e^2$ intensity spot size of the laser is about 300 $\mu$m in diameter, and the peak intensity is up to $5 \times 10^{14}$ W/cm$^2$.

The rise time of the laser pulse is about 150 ps. The electrons have a spot diameter of about 250 $\mu$m FWHM. Because the source of electrons is a 9.3 GHz RF linac, there is a micropulse structure with a 107 ps separation. The width of the micropulses was measured (by time resolving Čerenkov light) at the interaction point to be less than 10 ps, giving a peak current of more than 270 mA. Each micropulse contains about 17 million electrons.

The optical diagnostics are back- and forward-scattered CO$_2$ laser light (spectrally resolved but time integrated), collective Thomson scattering of a probe laser (frequency doubled output from a pulsed 1 pm laser), and frame-grabbed images of the plasma self-emission. The Thomson scattering has two configurations. The first has the probe beam focused to a point and $k$-matching set for $2k_0$ modes of the plasma, which are either electron plasma waves (“slow waves”) due to stimulated Raman/Compton scattering (SRS or SCS) or counterpropagating optical mixing (CPOM) or ion acoustic waves due to stimulated Brillouin scattering (SBS). Here, $k_0$ is the wave number of the 10.6 $\mu$m pump. In the second configuration, the probe beam has a line (cylindrical) focus and scattered light is collected at an angle corresponding modes with $k_\parallel = \Delta k$ where $k_\parallel$ is the component of $k$ parallel to $k_0$, and $\Delta k$ is the wave number difference between the two CO$_2$ laser pump beams. The signal in this second configuration is proportional to the level of the relativistic plasma wave (the “fast wave”). In both configurations, the scattered light is frequency and time resolved.

Figure 1 also shows a few electron trajectories. Both the accelerated and nonaccelerated electrons enter a
variable-field, imaging electron spectrometer. The nonaccelerated electrons are dumped onto low-density plastic. Lead shielding reduces the flux of background x rays reaching the electron detectors, thereby reducing the background or noise levels to a value negligibly small compared to the signal levels ultimately obtained. The accelerated electrons exit the vacuum through a 25 μm thick Mylar window, and are detected electronically by one or more silicon surface barrier detectors (SBD) or by a gaseous Čerenkov cell followed by a photomultiplier tube (PMT). The Čerenkov cell/PMT combination acts as a threshold detector producing no signal unless the electron velocity exceeds the speed of light in the working gas. This threshold is easily adjusted to fall in the range of 10–30 MeV. Electrons are also detected photographically by the tracks they leave in a cloud chamber. The SBD has a 300 μm copper window, which is “light tight” to soft x rays, but still “transparent” to energetic electrons. Various charge-sensitive preamplifiers and thick lead apertures are used to extend the dynamic range of the SBDs to over four orders of magnitude in electron flux, beginning with single-electron sensitivity.

In addition to these diagnostics for the accelerated electrons, there are two more electron diagnostics to determine the fate of those electrons which remain near 2 MeV. The first is a thin glass slide located at the beam dump. The strength of the bending magnet can be chosen so that the 2 MeV electrons enter the glass slide at the Čerenkov angle (about 47°) so that a portion of the Čerenkov light generated within the slide will be directed down the long axis of the slide. The end of the slide is directly imaged onto the photocathode of the same streak camera used for the slow-wave Thomson scattering. Thus, a record of the 2 MeV “transmission” of the plasma as a function of time can be obtained, and can be compared to the time history of the various plasma waves on a single shot. The time resolution for the Čerenkov signal is sufficient to distinguish the string of micropulses, but not enough to see changes within a single micropulse. The time resolution of the Thomson scattering at the same time is about 20 ps. In addition to quantifying the electron transmission, this diagnostic serves another purpose as well. While the timing of the linac macropulse and the CO₂ laser pulse is deterministic (by using an amplified version of the short, few hundred ps laser pulse to trigger the gun of the linac) the phase of the 9.3 GHz RF is random from shot to shot so that the timing of the micropulses is not deterministic. The consequence of this is that only a small fraction of the laser shots will produce the highest-energy electrons, since the lifetime of the peak plasma wave fields may be only 10 or 20 ps, as discussed later. However, the Čerenkov slide in the beam dump can be used to monitor the timing of the micropulses with respect to the plasma wave with an accuracy of about 20 ps, thus helping to interpret the accelerated electron data.

The second diagnostic for the 2 MeV electrons is a simple fluorescing screen viewed by a sensitive CCD camera. For these measurements, the beam dump is removed, the fluorescer is stretched across the exit aperture of the vacuum chamber, and the field in the bending magnet is lowered to direct the 2 MeV electrons to the fluorescer.

III. EXPERIMENTAL RESULTS

This paper is not meant to be a comprehensive review of all the experimental results that are reported elsewhere.³-⁵ It will, instead, focus on two particular measure-
of a normalized to the difference frequency. A white bar indicates the location of light from density fluctuations with a $k$ around $2/c$. Frequency shifts are second- and third-harmonic features from (a). (c) Lineouts of the Čerenkov light from the 2 MeV electrons (solid curve) recorded simultaneously and dashed curves, respectively. The signals at $2\Delta\omega$ and $3\Delta\omega$ are clearly visible. Figure 2(b) shows temporal lineouts through the streak record at $2\Delta\omega$ and $3\Delta\omega$ as the solid and dashed curves, respectively. The signals at $2\Delta\omega$ and $3\Delta\omega$ cannot be harmonics of a slow wave, because these would have a $k$ near $4k_0$ or $6k_0$, respectively, and therefore not visible in this $2k_0$ Thomson scattering setup. Instead, the signals must be either mode coupling of the second and third harmonics of the fast wave to an ion acoustic wave or mode coupling of the fundamental and second harmonic of the fast wave to a slow wave. Since $\Delta k/2k_0=0.015$, these mode-coupled waves have a $k$ that is still around $2k_0$, and will therefore be visible in this scattering setup. In either case, the signals observed at $2\omega$ and $3\omega$ are representative of the timing (if not the exact time history) of the fast wave. Lineouts from a later experimental run with the fast wave Thomson scattering configuration confirm that the plasma wave has a duration of 100–150 ps FWHM, as seen here.6

The simultaneously measured time history of the transmitted microbunches is shown in Fig. 2(c). The relative zero for the graphs in Figs. 2(b) and 2(c) was determined to within 50 ps by evacuating the vacuum chamber and watching the effect of the laser itself on the 2 MeV electrons. The transverse kick from the laser ponderomotive force is sufficient to just disturb the Čerenkov signal, from which we can determine the timing of the peak of the laser with respect to the linac.7 We then fire into the plasma and assume that the low level Compton scattering (seen along with the bright $\Delta\omega$ features) is centered around the peak of the CO$_2$ laser pulse. Thus the Thomson scattering signal and electron signal can both be referenced to the peak of the CO$_2$ pulse to within the 50 ps shot-to-shot jitter of the triggering circuits.

Figure 2(c) shows that the 2 MeV electrons have a sudden and deep drop in transmission, coincident with the onset of the mode coupled feature, which we interpret as being indicative of the presence of a fast wave. We also see that the transmission begins to recover after $\approx 1$ ns, but then drops again later in the pulse. These are features of shots where a large-amplitude plasma wave was excited, as evidenced by the detection of high-energy electrons on the SBD detectors. Other laser shots with weak plasma waves show a much smaller effect, i.e., a smaller drop and a faster recovery with no late-time recurrence of the drop in transmission, while single-frequency shots show an even smaller effect, with typically a 300–400 ps full-recovery time.7

The Čerenkov signal could be reduced from the no-laser case for a variety of reasons. First, the energy of the electrons can change so that they no longer strike the slide or strike it at the wrong angle for producing Čerenkov light. Second, the number of 2 MeV electrons reaching the slide can be reduced if the electrons are scattered into an angle larger than the acceptance angle of the measuring device. Third, the number of 2 MeV electrons reaching the slide can be reduced if the entire beam is deflected enough to miss the aperture of the Čerenkov measurement. We believe that all three effects are happening during beat wave shots, as discussed below.

Accelerated (and decelerated) electrons are observed in the experiment, and this must deplete the 2 MeV electrons that become trapped in the accelerating and focusing phases of the wave. The electrons that are not trapped in the accelerating and focusing phases will likely leave the wave at angles larger than their initial angle due to the large magnitude of the radial fields. To give us a rough idea of the numbers of 2 MeV electrons exiting a model plasma wave at various angles, we numerically pushed electrons (with the experimental parameters) through the 2-D fluid model of the longitudinal and radial fields by Fedele et al.,8 and looked at the output angle versus output energy. For a laser beam with our experimental axial and radial scale lengths, driving a peak wave amplitude of 30%, the reduction of electrons at $2\pm 0.15$ MeV into the experimental collection angle ($\pm 3^\circ$) is a factor of about 6. Thus, the

**FIG. 2.** (a) A streak image of the spectrally resolved Thomson scattered light from density fluctuations with a $k$ around $2/c$. Frequency shifts are normalized to the difference frequency. A white bar indicates the location of a 10X attenuator over a portion of the spectrum. (b) Lineouts of the second- and third-harmonic features from (a). (c) Lineouts of the Čerenkov light from the 2 MeV electrons (solid curve) recorded simultaneously with the streak data of (a) and (b), and of the Čerenkov light from a reference shot in which the CO$_2$ laser was not fired (dashed curve).
sudden drop in transmitted electrons is the expected consequence of having a coherent, accelerating electron plasma wave with a small diameter. On the other hand, Thomson scattering indicated that this coherent plasma wave disappears after 100 ps or so while the reduction in transmission persists on ns time scales. How can the lack of 2 MeV electrons at later times be explained? We can get more insight into this paradox by viewing the actual profile of electrons striking the Čerenkov slide.

As discussed in Sec. II, to view the electron beam profile the electron spectrometer magnetic field is reduced allowing the 2 MeV electrons to strike a fluorescer placed at the exit plane of the vacuum chamber. Figure 3 shows various beam profiles for various timings of the peak of the linac macropulse with respect to the peak of the CO$_2$ laser pulse. Recalling that the macropulse has a 1.5 ns FWHM, a timing of negative 1 ns puts most of the electrons through the interaction point before the CO$_2$ pulse arrives. This is shown in Fig. 3(a). This profile is not a perfect, round spot due to aberrations in the imaging properties of the double-focusing sector magnet and because the fluorescer is not in the exact focal plane of the magnet. The acceptance angle of the magnet is slightly larger than the focusing angle of the electrons, so that this image represents nearly all the electrons injected into the experiment. With the relative timing set to 0 ns, the bright central core of the image is reduced by about a factor of 2. This is similar to the situation shown in Fig. 2(c), where somewhat more than half of the electron signal is missing. In Fig. 3(c), the timing is +1 ns and the bright core is almost completely gone, leaving a diffuse background. Finally, Fig. 3(d) shows the profile at a time of 5 ns after the peak of the CO$_2$ pulse. This is well after the transmission has made its first, 0.5–1 ns time-scale recovery but not its second. We see that the image again has a bright core, but that it is offset as if the beam were deflected vertically. The direction and magnitude of this vertical offset appears to be random from shot to shot.

Our interpretation of these observations is the following. The transmission drop in the Čerenkov data between 0
and 0.5–1 ns is due to diffuse electron scattering into an angle larger than the ±3° acceptance angle of the measurement. This is consistent with the diffuse image in Fig. 3(c). The drop in transmission in the Čerenkov data beyond 1 or 2 ns is more due to the electron beam being deflected as a whole, and either missing the aperture of the magnet or striking the Čerenkov slide at a nonoptimum angle, as is suggested by the offset beam image seen in Fig. 3(d). We believe that this later deflection is due to strong magnetic fields generated in the plasma from a Weibel-like instability as the plasma wave damps. We will be looking into this possibility with improved electron imaging capabilities in a later experiment.

To confirm that the diffuse images may indeed be the result of diffuse scattering of the electrons at the interaction point, we inserted a variety of thin foils into the beam path at the interaction point and observed the resulting images. We found that we could reproduce images similar to Fig. 3(c) with a 15 μm thick Mylar foil, indicating that the RMS scattering angle in Fig. 3(c) is about ±10°. We should note that the placement of scattering foils at the interaction point did not lead to false signals on any of the accelerated electron detectors. Thus, we have learned that the long, 0.5–1 ns recovery time in the electron transmission can be accounted for if there exists a low level plasma wave with sufficient radial fields to scatter the incident electrons into an angle of 3°–10°.

B. Cumulative energy spectrum

The flux of electrons at various energies was quantitatively measured by the silicon surface-barrier detectors. These detectors have a depletion depth (active thickness) of about 1 mm. The range of an electron in silicon is greater than 1 mm for energies greater than 0.7 MeV, and therefore, for all the energies measured in this experiment (1–30 MeV), the electron will penetrate through the detector. Thus, the signal measured will be \( (dE/dx)_{\text{ion}} \times dx \times C/eV \times mV/C \), where \( (dE/dx)_{\text{ion}} \) is the energy loss rate (eV per mm) due to ionization in the silicon, \( dx \) is the thickness of the depletion region (1 mm), \( C/eV \) is the number of electron/hole pairs generated per eV of energy loss expressed in coulombs per eV, and \( mV/C \) is the charge-sensitive preamplifier gain expressed in mV per coulomb. For the high-gain preamp, this gives 20 mV/electron at 2 MeV. The energy dependence of the signal per electron comes through the energy dependence of \( (dE/dx)_{\text{ion}} \). This is flat to within 20% from 1 to 30 MeV. By knowing the energy location of the detector on a given shot, one can convert the mV signal into detected electrons. Given the aperture of the detector and the local dispersion of the magnet (MeV/mm), one can convert the measurement to electrons MeV. The vertical spread of the accelerated beam is folded in by assuming that the angle varies as \( 1/\gamma \), the Lorentz factor at that energy.

The resulting energy spectrum obtained over many laser shots is shown in Fig. 4. We see an experimental spectrum that falls off rapidly with energy on both sides away from 2 MeV and extending out to about 30 MeV on the high-energy side. We should emphasize that the experimental parameters were not held constant for all these shots. There were variations in density (typically only a few percent around the resonant value), laser energy, line ratio, and, most importantly, timing of the micropulses, with respect to the peak fields of the accelerating electron plasma wave. Thus, the highest points are more representative of the idealized conditions one would assume in a model.

IV. MODEL OF PLASMA FORMATION AND WAVE GROWTH

The simplest model of the nonlinear growth and saturation of the plasma wave was given by Rosenbluth and Lui, and extended for arbitrary densities and laser intensities by Tang et al. These were single-point treatments in Lagrangian coordinates. The 2-D problem was studied by Fedale et al., who solved the linear fluid equations in cylindrical coordinates. However, the solution essentially describes a thin slice of plasma and does not include longitudinal variations due to laser focusing, finite transit time of the plasma, or laser rise time. Other effects such as ionization dynamics, and, most importantly, relativistic saturation, were also neglected. Stated this way, a full understanding of the problem is actually beyond simple numerical computations and is best suited for particle-in-cell (PIC) simulations.

In beat wave simulations using the 2-D PIC code done by Mori, strong radial and axial convolutions in the wave amplitude were observed at late times due to the intensity-dependent saturation time across the laser spot and the axial transit time of the simulation box. He observed that for laser spot sizes much larger than \( c/\omega_p \), the collisionless skin depth, and until the radial structure scale length became comparable to \( c/\omega_p \), that the time behavior at each point in the simulation box was roughly what you would expect for a Lagrangian oscillator located at that point. We therefore believe that a great deal of insight into the spatial–temporal dynamics can be gleaned by solving the 1-D Tang et al. formulism numerically on a 1-D or a 2-D grid, treating each point independently. The density buildup at each point can be calculated from the
tunneling ionization rate and any density reduction from hydrodynamic expansion of the plasma due to thermal or ponderomotive forces can be included, so that a rather complete modeling of the experiment can be done in this way.

For these calculations, a model laser pulse is assumed, which closely matches the experimental laser parameters. The rise time (fall time) is 150 (200) ps and the spatial envelope is assumed to be a quasi-Gaussian with a 1/e radius (for the field) of 150 μm and a Rayleigh length of 0.8 cm. The time history at each point on the grid differs from any other point only by the time at which the laser reaches that point and the peak intensity ultimately obtained at that point. Thus, a table of numbers representing the time histories (rows) at various peak intensities (columns) of the beat wave amplitude can be used as a look-up table to look up the amplitude at any (r,z) point and at any time. Thus, one can construct 2-D images of the wave amplitude at various times or find the wave amplitude in the frame of a comoving electron to work out the energy gain of that electron.

The time histories of the plasma density and the plasma wave amplitude are solved numerically by integrating the tunnel ionization rate equation simultaneously with the ponderomotively driven plasma wave equation from Tang et al. The time arrays for the instantaneous plasma wave amplitude are analyzed in a Hilbert transform subroutine to extract the envelope amplitude, the instantaneous plasma wave frequency, and the instantaneous shift in phase from the beat wave driver. In this way, we construct look-up tables for the plasma density, the laser field, the plasma wave amplitude (envelope), the plasma wave frequency, and the phase shift between the plasma wave and the ponderomotive beat wave. These tables are generated for various initial neutral densities and further analysis, e.g., working out spatial profiles or electron energy gain versus time, is done by post-processing on these tables.

An example of the output of this model is shown in Fig. 5. Figures 5(a) and 5(b) are two frames from the evolution of the laser intensity, plasma density, plasma wave amplitude, and plasma wave phase taken at t = 103 and 168 ps, respectively. The laser propagates from left to right and is focused at the center of the box at x = z = 0. The time t = 0 is when the laser pulse is just arriving at z = 0, so that the first frame is 47 ps before and the second frame is 18 ps after—the peak of the laser pulse. Relativistic saturation has just occurred in the first frame near the center of the box. Figures 5(c) and 5(d) show the on-axis lineouts corresponding to the two frames, i.e., the 1-D representation of the model. In this calculation, the neutral pressure was set to 10% above the “resonant” pressure, an initial offset that experimentally seems to be near optimum. Upon full ionization, the plasma was allowed to drop in density linearly in time at a rate of 0.15 n_{res} per 100 ps, where n_{res} is the theoretical resonant density (n_{res} = 9.4 × 10^{15} cm^{-3}). This is the experimentally measured rate at which the plasma frequency typically drops, as measured by the time-resolved frequency shifts in the fast wave Thomson scattering setup. The plots of plasma density show the tunneling ionization front moving to the right and the linear drop in plasma density versus time, showing up as a spatial ramp in density as one moves back to the left. The plasma wave amplitude saturates at n/n_{0} = 33% at t = 103 ps.

The phase of the wave relative to the beat wave driver is somewhat complicated due to the plasma expansion. We can see in the lineouts in Fig. 5(c) that the phase variation with position is least rapidly varying, where there is strong coupling of the pump to the plasma wave—that is, for z values, such that the density is within the resonance of the harmonic oscillator. At t = 168 ps, the wave amplitude near z = 0 has not changed much, but, since the density at that point has dropped outside the resonance, the wave and driver have decoupled, and the phase has evolved according to the difference between the driver and local plasma frequency. (All phases are displayed between 0 and 2π, but are actually continuously varying. A finer calculation grid would show a phase jump at 0 or 2π.)

V. DISCUSSION

The maximum electron energy gain in the plasma wave is given by ∫ E_{x} dy dl and so is essentially proportional to the peak wave amplitude and the effective length of the wave. If one looks at the lineouts of the wave amplitude in Figs. 5(c) and 5(d) for t = 103 and 168 ps, respectively, one might guess that the energy gain should be larger at the later time, since the peak amplitude has not changed much, but the FWHM length has almost doubled from 1.0 to 1.8 cm. However, the phase variation in the frame of the accelerating electron needs to be considered. We have calculated the maximum possible energy gain versus time from our model by pushing many groups of electrons through the wave at start times separated by 10 ps. Each group consisted of ten electrons spread out over one plasma period about 1 ps). These groups were accelerated in one dimension along the z axis, according to the local amplitude and phase determined as before, and the maximum energy obtained was recorded as a function of the time at which that group crosses z = 0. This is plotted as the solid curve in Fig. 6. We see a narrow peak in energy gain versus time at about 103 ps. Thus, the most optimally placed electron is one that arrives at the location of the peak wave amplitude, just as the wave is saturating. Although the wave amplitude is still large 65 ps later, the effective phase velocity of the wave changes rapidly and puts the electrons into decelerating phases. We can therefore take the effective acceleration length as 1 cm. Combining this with the experimentally measured energy gain of 28 MeV implies a peak accelerating gradient of 2.8 GeV/m and a peak wave amplitude of 28%. The energy gain in the idealized 1-D model is 57 MeV, which is about a factor of 2 greater than the largest gain yet observed in the experiment. However, since the energy gain is highly sensitive to phase coherence along the electron trajectory and real electron trajectories can never be strictly 1-D, additional 3-D effects, such as radial fields, neglected in this model, may limit the ultimate energy gain.

We have yet to address the paradox proposed in Sec. III A in which the Thomson scattering lifetime was measured to be on the order of $100 \text{ ps}$, while the apparent scattering of the injected 2 MeV electrons persisted for about 1 ns. Can the dephasing we observe in the model that limits the time over which electrons can gain energy, also account for a narrow time window for observing collective Thomson scattering? To address this question, we note that the Thomson scattered power per unit frequency interval is proportional to the spectral density function $S(k, \omega)$, which itself is related to the square of the Fourier transform of the spatial (and temporal) distribution of the electron density perturbations. In our experiments, the green probe beam has a 6 mm line focus centered at $z=0$. We therefore perform a 2-D Fourier transform on the middle 6 mm of the computation box, using both the amplitude and phase of the wave, as given in, for example, in Figs. 5(a) and 5(b) for $t=103$ and 168 ps. We are neglecting any changes to the amplitude or phase for the transit time of the probe beam, but this is a good approximation for the 90° angle of incidence (the transit time is 2 ps across the box). The result is shown as the dashed curve in Fig. 6. Here, we plot the square of the amplitude of the wave at $k_\parallel = k_\perp = \Delta k$, which is the location in $k$ space for which the Thomson scattering is $k$ matched. This model predicts that the scattered power will peak at $t \approx 100 \text{ ps}$—again, just

FIG. 5. A 2-D model of the laser field, electron density, plasma wave amplitude, and phase relative to the driver obtained by applying the Lagrangian oscillator model with tunneling ionization to each point on a 2-D Cartesian grid. Spatial variation are shown for (a) $t=103 \text{ ps}$ and (b) $t=168 \text{ ps}$. Shown in (c) and (d) are the 1-D analogs to (a) and (b), which are essentially lineouts down the axis.
when the wave is saturating—and fall off with a FWHM of about 50 ps. Thus, the phase aberrations of the wave may also limit the lifetime of the collective Thomson scattering.

We have not solved for any radial fields in this model, and so we cannot push electrons through and look at the lifetime of the scattering of electrons around 2 MeV. However, this model suggests that those radial fields would persist for a long time. This is because the damping on the wave is very weak. Collisional damping is negligible on ns time scales and damping onto slow waves through mode coupling also appears to be a slow process, with damping time scales predicted to be longer than a ns. Landau damping is also unimportant, as the numbers of injected particles is very small. We therefore expect that the large fields of the plasma wave do persist and scatter the injected electrons into large angles (> 3°), while perhaps providing some small energy gain. Clearly, this problem calls for a more accurate 2-D simulation. We are in the process of modifying the 2-D PIC code WAVE to simulate the problem on realistic experimental time scales. This requires that the ponderomotive force be introduced explicitly into the code (replacing the two high-frequency pump waves) so that WAVE can solve for the motion of the electrons and ions on > 100 ps time scales.

VI. SUMMARY

In summary, using a two-frequency CO₂ laser, we have excited large-amplitude electron plasma waves through the plasma beat wave process. Electrons with an energy of 2 MeV were injected into the plasma wave, and the observed output energy spectrum extended to 30 MeV on the high-energy side and 1 MeV on the low-energy side. The time history of the wave fields was studied through both Thomson scattering and through the perturbation to the beam of electrons at 2 MeV, with each diagnostic implying a different time scale for the lifetime of the wave fields. The Thomson scattering shows a short-lived wave (≈ 100 ps), while the perturbation of the 2 MeV electrons continues as a diffuse scattering for 0.5–1 ns after the Thomson scattering terminates, which then changes to a whole-beam deflection for at least the next 5 ns. A simple model suggests that the wave coherence may be short lived, while the wave fields themselves persist for a longer time. This can lead to a narrow time window for electron acceleration and for Thomson scattering, while the scattering of the 2 MeV electrons can persist on longer time scales. The whole beam deflection may be due to magnetic fields from a Weibel instability caused by the relaxation the anisotropic electron distribution function resulting from the eventual damping of the plasma wave. This will be a topic of further study.

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2See, for example, the Special Issue on Plasma-Based High Energy Accelerators, IEEE Transactions on Plasma Science PS-15 (1987).