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COLLECTIVE BEHAVIOR IN RECENT LASER-PLASMA EXPERIMENTS*

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ABSTRACT

The absorption of 1.06μ light focused on small C_6H_6 discs has been measured in the intensity range of $10^{15} - 10^{17}$ W/cm². The data confirms the importance of collective plasma effects in determining the absorption of intense light. The measured absorption efficiencies are in the range of 30-40%, and the scattered light has a polarization dependence. The measured absorptions are shown to be inconsistent with classical inverse bremsstrahlung. Both the magnitude of the absorption and the observed polarization dependence of the scattered light are shown to be consistent with recent calculations of light absorption via collective processes. Comparisons are made, and improved models of the light absorption are discussed with reference to the data. In addition, the heated electron energies deduced from the X-ray data are consistent with those expected via collective processes.

INTRODUCTION

The interaction of intense laser light with plasmas is a problem rich in nonlinear effects. For example, for Nd laser light with an intensity of 10^{16} W/cm², the energy of oscillation of electrons in the free-space electric field of the light is ~ 4 keV, a value comparable to the thermal energy of the plasma.

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Under such conditions, classical inverse bremsstrahlung (collisional) absorption is quite weak, but many collective mechanisms are operative. Some of these processes enhance the absorption by coupling the incident light into plasma waves which are then damped. Such processes include resonance absorption,¹⁻⁴ parametric instabilities both near the critical density^{5,6} and near one-fourth the critical density,^{7,8} and quasi-resonant coupling on ion density fluctuations^{9,10} driven up by subsidiary processes such as by two-stream instabilities in the plasma expansion. On the other hand, there is in general also stimulated scattering of the incident light in the plasma with density less than critical via the Raman and Brillouin instabilities.¹¹⁻¹³

All these processes are operative in a highly inhomogeneous, expanding plasma and their evolution is strongly modified by the reaction of the turbulence back on the density and velocity gradients.^{3,4,11} In addition, the absorbed energy fluxes can be far above those for which classical theory is valid, and the energy transport can be determined by heat-transport instabilities¹⁴ or self-generated DC magnetic fields.¹⁵⁻¹⁸ Although the understanding of the competition between these many effects is of course far from complete, it is encouraging that correlations with experiments are increasing. It is the purpose of this paper to discuss a number of these correlations by relating recent calculations with some recent experimental results. We will show that there is considerable evidence that collective plasma processes play an important role in determining the absorption of intense light.

MEASUREMENTS OF LIGHT ABSORPTION

There have been many previously reported experiments¹⁹⁻²¹ on the irradiation of targets of various types with intense laser light. In the experiments discussed here, small CaH_2 discs (150 μ diameter by 10 μ thick) were irradiated with 1.06 μ radiation using the Janus laser facility.³² The energies on target ranged from 5-15 J in pulses with a full-width at half-intensity ranging from 60-150 ps. The light was focused to a target spot with diameters of approximately 90, 30 and 10 μ . The resulting irradiation intensities ranged from 10^{15} - 10^{17} W/cm². Extensive diagnostics were taken, including measurements of optical energy balance, the angular distribution and time-dependence of the reflected light, the X-ray energy spectrum from $\sim$ 1-100 keV, the ion expansion energy, and the spatial distribution of the X-ray emission. Details of these experiments and numerical simulations of them using a 2-D fluid-dynamics code with multi-group electron transport (LASNEX)³³ are discussed in Ref. 17 and 18. We will here concentrate on the role of collective plasma effects on the light absorption, as evidenced by the data.

Fig. 1

The fraction of the incident laser light absorbed as a function of incident intensity I_L . The crosses denote measurements (with estimated error bars) determined by optical energy balance using an array of photodiodes distributed in angle around the target augmented by calorimeters to measure the light reflected and transmitted through the lenses. The points denote measurements in which the photodiodes are replaced by a box calorimeter enclosing the target.

Fig. 1 shows the measured absorption as a function of intensity. As indicated, on most shots the absorption was determined from measurements of the scattered light using ~ 22 calibrated photo-diodes which were distributed in angle around the target (and calorimeters for the light scattered into the lens). The absorption measured by this technique ranges from $.27 \pm .09$ at a nominal intensity of 2×10^{15} W/cm² to $.41 \pm .08$ at $3-5 \times 10^{16}$ W/cm². On a few shots, the photo-diodes were replaced with a calorimeter which enclosed the target. The absorption determined by this alternative technique ranges from .35 to .39 for the above intensities. The magnitudes of the absorption given by the different techniques are in reasonable agreement, but clearly the trend for the absorption to increase with intensity must be considered tentative; i.e. within the uncertainty of the measurements.

Fig. 2 shows the fraction of the incident light energy reflected back through the lens. This fraction decreases from $\sim 35\%$ to $\sim 15\%$ as the intensity increases in the range of 10^{15} - 10^{17} W/cm². Fig. 3 shows some scattered light distributions averaged over the photodiodes in and out of the plane of polarization. Note that the scattered light peaks in the backwards direction ($\theta = 180^\circ$) and that very little light is refracting around the target. As the incident intensity increases, the scattered light distribution remains peaked in the backwards direction but becomes somewhat broader. Lastly, these measurements show that the scattered light has a polarization dependence. As shown in Fig. 4, photodiodes perpendicular to the plane of polarization collect on the average more reflected light than do photodiodes in the plane of polarization. The asymmetry ranges from a factor of ~ 1.2 to a factor of ~ 3 depending on scattering angle and intensity.

Fig. 2

The fraction of the incident light energy reflected through the focusing lens as a function of incident intensity. The measurements are made with a calorimeter, and the estimated error bars are shown.

Fig. 3

The azimuthally averaged fraction of the incident laser light energy scattered per unit solid angle as a function of the scattering angle θ . Direct backscatter corresponds to $\theta = 180^\circ$. The circles denote the average value measured with photodiodes in experiments with a 30 μ diameter target spot, and the crosses denote the average value in experiments with a 90 μ diameter spot.

Fig. 4

The ratio of the fraction of the light energy scattered per unit solid angle perpendicular to the plane of polarization to the

fraction scattered parallel to the plane as a function of the cosine of the scattering angle θ . The circles denote averages over several experiments with intensities in the range $1.9 \pm .2 \times 10^{15}$ W/cm², and the squares denote averages over several experiments with intensities in the range $2.0 \pm .4 \times 10^{16}$ W/cm². The measurements are made with photodiodes.

As we will discuss, the measured absorption (and its polarization dependence) does not agree with inverse bremsstrahlung absorption, but is in semi-quantitative agreement with that expected via collective processes in the plasma. Let us first address the inverse bremsstrahlung and show that it is rather small for the higher intensities investigated. Then we will relate recent calculations³ of light absorption via collective processes to these experiments.

CLASSICAL INVERSE BREMSSTRAHLUNG

It is appropriate to first examine the classical collisional absorption. The absorbed intensity in a linear density profile is then^{3,4}

$$I_{\text{ABS}} = I_{\text{INC}} [1 - \exp(-\frac{32}{15} k_0 L \frac{\nu_{\text{cr}}}{\omega_0})] .$$

Here I_{INC} is the incident intensity, k_0 the free-space wave number, L the gradient length, and ω_0 the light frequency. ν_{cr} is the electron-ion collision frequency evaluated at the critical density:

$$\nu_{\text{cr}} = 3.4 \langle Z \rangle \ln \Lambda \frac{n_{\text{cr}} e^4}{m_e^{1/2} \theta_e^{3/2}} ,$$

where n_{cr} is the critical density, θ_e the electron temperature, e the electron charge, m_e the electron mass, $\langle Z \rangle$ the average charge state of the ions, and Λ the usual ratio of the maximum to the minimum impact parameter. It is sufficient for these estimates to express the absorbed energy flux as

$$I_{\text{ABS}} = f n_{\text{cr}} \theta_e \nu_e ,$$

that is, as a factor f times the free-streaming value. Note that if we take $f = 1$, we are finding the minimum temperature consistent with a given flux of absorbed energy, and hence the maximum clas-

sical absorption.³⁵

Putting these equations together gives $1 - \epsilon = e^{-B/\epsilon}$, where ϵ is the absorption efficiency ($I_{\text{ABS}}/I_{\text{INC}}$) and

$$B = 2 \times 10^{10} \frac{\langle Z \rangle \ln \Lambda f}{I_{\text{INC}}} k_0 L.$$

I_{INC} is in W/cm^2 . If $B \gg 1$, then $\epsilon \approx 1$. If $B \ll 1$, then $\epsilon = \sqrt{B}$.

To illustrate the numbers, consider some parameters characteristic of these short-pulse experiments: $k_0 L \approx 10^2$, $\langle Z \rangle = 3$, $\ln \Lambda = 5$ and $f = 1$ (to find the maximum inverse bremsstrahlung). Then

$$B = \frac{3 \times 10^{13} \text{ W}/\text{cm}^2}{I_{\text{INC}}}.$$

For $I_{\text{INC}} = 10^{16} \text{ W}/\text{cm}^2$, $B \ll 1$ and $\epsilon \approx .06$. Note also that in this regime, $\epsilon \propto I_{\text{INC}}^{-1/2}$.

As a matter of fact, there are a number of reasons why even this small value is an over-estimate of the inverse bremsstrahlung absorption. First, as already noted, the condition $f = 1$ assumes essentially the most efficient energy transport possible (and hence the lowest electron temperature consistent with the collisional absorption). As discussed in Ref. 17 and 18, the transport appears to occur at a rate much less than this free-streaming value, i.e. $f \ll 1$. Furthermore, we have neglected in this simple calculation several nonlinear effects which reduce the collisional absorption. These nonlinear effects are the intensity corrections to the collision frequency and the local steepening of the density profile near the critical density due to ponderomotive forces. Our principal point is that neither the magnitude nor the intensity dependence of the measured absorption is consistent with collisional absorption.

COLLECTIVE ABSORPTION PROCESSES

On the other hand, the measured absorptions are consistent with those found in recent calculations of laser light absorption via the generation of electron plasma waves. To illustrate this point, we first refer to some two-dimensional simulations³ of laser light absorption in inhomogeneous plasma which were carried out with a 2-D particle code. In these simulations uniform plane

waves were propagated from vacuum into an inhomogeneous plasma slab, and the absorption was computed as a function of angle of incidence.

Some representative results are shown in Fig. 5, which is a plot of absorption efficiency versus angle of incidence. For these results the plasma slab had an initially linear density profile with scale length $k_0 L = 10$ and an initial electron temperature of 4 keV. The incident light had an intensity of 10^{16} W/cm^2 (1.06 μ radiation) and was p-polarized; i.e., had its electric vector in the plane of incidence. Note that the fractional absorption computed in these short-term simulations (the points in Fig. 5) ranges from $\approx 15\%$ at normal incidence to $\approx 60\%$ at an angle of incidence to $\approx 20^\circ$.

Fig. 5

The absorption efficiency as a function of angle of incidence θ . The points are the results of 2-D particle simulations¹ of p-polarized light incident on a plasma slab. The dashed line is an estimate, based on this data, of the net absorption efficiency for light incident with an equal mix of polarizations. The simulation results are for a plane wave with an intensity of 10^{16} W/cm^2 (1.06 μ light) incident on a plasma slab with an initial density profile varying linearly from 0 to $1.7 n_{cr}$ in a distance of 13 vacuum wavelengths. These results should only be considered as representative.

These absorptions are consistent with those expected in the nonlinearly steepened density profiles which are self-consistently generated in the simulations by strong ponderomotive forces near the critical density. An important effect of this strong local steepening is to reduce the efficiency of parametric instabilities and make resonance absorption effective over a rather broad range of angles. A few estimates will suffice to make these points clearer. The absorption efficiency due to parametric instabilities near the critical density in the steep profile is $\epsilon_{pi} \sim 2\nu^*/\omega_0 k_0 l$, where ν^* is the effective damping rate and l is the size of the region of plasma accessible to the instability excitation. At high intensities $\nu^* \sim .05 \omega_0$, and l is $\sim$ few c/ω_0 in the nonlinearly steepened profile. Hence instability generation in the steep profile is expected to provide at most a rather modest absorption of order 10-20%, consistent with the simulations.

The absorption efficiency due to resonance absorption of p-polarized light can be estimated by computing the energy damped by the resonantly-excited electron plasma waves. For a locally $1/3$ linear profile, the result is $C_{ra} \sim |\phi(\tau)|^2/2$. Here $\tau = (k_0 L) \sin \theta$, θ is the angle of incidence, and $\phi(\tau)$ is the well-known resonance function which has a maximum value of ≈ 1 at $\tau = .8$. Of course, the details of both the resonance absorption and the instability generation depend on the actual shape of the nonlinearly steepened profile, but the qualitative features of the calculated absorption are consistent with these simple estimates.

In the simulations, the incident light is p-polarized. In experiments, generally only part of the incident light ($\sim 1/2$) has this polarization. The dashed line in Fig. 5 is our estimate, based on this simulation data, of the average absorption which results when one assumes an equal mix of polarizations for the incident light. Hence this estimate takes into account the fact that resonance absorption is only operative on one polarization. The resulting prediction is that the absorption via plasma wave generation is (1) polarization-dependent and (2) has a net efficiency of ~ 15 -30% as a function of angle of incidence.

Both these results compare favorably with the data. Since the light was focused with an $f/1$ lens, the angle of incidence averaged over energy on a plane target is $\approx 15^\circ$. We then estimate from Fig. 5 about 25% net absorption due to plasma wave generation. Since p-polarized light is absorbed more efficiently than is s-polarized light (the electric vector out of the plane of incidence), we also estimate that diodes out of the plane of polarization collect ~ 1.5 times as much reflected (nonabsorbed) light as do diodes in the perpendicular plane. Note from Fig. 4 that the observed asymmetry is not large (1.1-1.5) for most of the scattered light, which is reflected into angles within ~ 30 - 40° of the direction of incidence. The principal point of these

comparisons is simply that the observations are consistent with light absorption via collective mechanisms.

IMPROVED MODELS OF THE LIGHT ABSORPTION

Although these correlations are sufficient to establish the importance of collective mechanisms, there are still some considerable uncertainties in our quantitative understanding of light absorption via such processes. Short-term simulations of the simple model discussed above are too ideal to predict in detail the magnitude of the absorption and its dependences on angle of incidence and intensity. The magnitude of the measured absorption ranges up to nearly twice as large as that estimated on the basis of Fig. 5. The angular spread of the scattered light is larger than the spread due to reflection from a plane critical surface, and the polarization asymmetry of this scattered light is sometimes as large as a factor of 3 for large scattering angles. Time-dependencies in the back-reflected light are observed in the higher intensity experiments ($I > 10^{16}$ W/cm²). Finally, the absorption efficiencies measured in some more recent experiments using a $f/2.5$ lens are approximately the same as those measured with the $f/1$ lens. This suggests a weaker dependence of the absorption on angle of incidence than is apparent from Fig. 5.

Improved models of the light absorption must include a number of additional effects not included in the simple model. Many of these effects both enhance the absorption and in particular reduce its sensitivity to angle of incidence. First, if the transport is severely limited, the plasma near the critical density surface can be at an extremely high temperature. This would reduce the region over which the profile is strongly steepened and allow for more instability-generated absorption. In addition, the model discussed above treated a uniform plane wave incident on an infinite plasma slab. In practice, the spatial intensity profile of a light beam can be (and is) strongly modulated on spatial scales of 10-100 μ . Such intensity modulations can crater the critical density surface either directly via variations in the ponderomotive pressure or indirectly by generating $\nabla n \times \nabla \phi$ magnetic fields which inhibit the transport and modify the fluid dynamics. Recent calculations show that a cratered critical surface can also be produced in other ways, for example, by a pinch-type instability¹⁶ of light which is initially self-trapped in channels along the surface. Clearly a strongly rippled surface both reduces the sensitivity of the absorption to angles of incidence and broadens the angular distribution of the scattered light.

In addition to these relatively long wavelength density variations (compared to λ_0 , the laser light wavelength), there

are also additional sources of short wavelength ion turbulence. For example, electron-ion instabilities induced by heat transport to higher densities¹⁴ and/or ion-ion instabilities driven by the expansion of different ion species can produce ion density fluctuations in the underdense plasma with a wavelength of $\sim 5 - 10 \lambda_{pe}$ (electron Debye lengths). These fluctuations enhance the absorption by coupling the light into electron plasma waves. The physics of the coupling is similar to that in absorption via parametric instabilities near the critical density, but now the ion fluctuations are generated by subsidiary processes rather than directly by the electric vector of the light. One important feature is that this mechanism operates over a broader range of densities and is less sensitive to a local profile steepening than are parametric instabilities.

It is instructive to give some simple estimates¹⁵ of this absorption for these experiments. The effective collision frequency due to the ion turbulence is

$$\frac{\nu_{eff}}{\omega_0} \sim \frac{1}{4} \left| \frac{\Delta n}{n_p} \right|^2 \frac{n_{cr}}{n_p} \left\langle \frac{1}{\epsilon(k, \omega_0)} \right\rangle$$

where Δn is the mean amplitude of the fluctuations which are assumed to be isotropic in angle, n_p the local plasma density, n_{cr} the critical density, ω_0 the light frequency, and ϵ the plasma dielectric function evaluated at ω_0 and at the wave number k of the fluctuations. The brackets denote an average over a range of wave numbers. To illustrate the numbers, take a range of wave numbers uniformly distributed from $k\lambda_{pe} = .5$ to $k\lambda_{pe} = 1$, use the well-known dielectric function for a plasma with a Maxwellian distribution of velocities, and estimate the mean fluctuation level as $\Delta n = .1 n_p$. This latter estimate is based on the amplitude of an ion wave at which ion trapping onsets in a plasma with $U_0 = 10 U_i$ (U_i is the ion temperature). Then near the critical density $\nu_{eff} \approx 1/800 \omega_0$. In a linear density profile with $k_{\perp} L = 10^2$, the estimated absorption is $\approx 20\%$, which is competitive with the net absorption estimated from Fig. 5. It is also interesting that this mechanism produces relatively soft heated electron distributions. The largest uncertainty at present is determining when such turbulence is actually present in the underdense plasma.

It is appropriate to briefly mention a few other collective mechanisms for absorption. For example, self-generated magnetic fields can introduce resonance absorption of even normally incident light.³⁷ In this case, the oscillation velocity of electrons in the electric field of the incident light resonantly drives electrostatic waves near the critical density via the Lorentz

force. Recent calculations indicate that this mechanism can introduce a modest amount of absorption ($\sim 10\%$) if the size of the self-generated magnetic fields is of order 5 MG. In addition, there is also the two-plasmon decay instability which operates near one-fourth the critical density. However, present calculations indicate that this instability is also strongly limited by a local profile steepening. Clearly a quantitative understanding of the many different absorption mechanisms is far from complete.

STIMULATED SCATTERING

More quantitative calculations of the absorption must also include an improved understanding of the role of stimulated scattering in the plasma with density less than critical. The ideal model referred to for the "zeroth-order" estimates included a rather small region of underdense plasma (of size ≤ 5 free-space laser light wavelengths) and so can underestimate this scattering. Several features of the data indicate that stimulated scattering, particularly via the Brillouin instability, is playing some role in these experiments. For example, stimulated side-scattering preferentially scatters light out of the plane of polarization and so also leads to a polarization dependence to the scattered light. Hence this scattering can account for the large asymmetry with respect to the plane of polarization observed in the light scattered at large angles. However, the light scattered into such angles is a small fraction of the incident energy.

In these experiments the intensity is increased by reducing the diameter of the target spot. Hence the size of the underdense plasma decreases as the intensity increases. For the experiments with $I \geq 10^{16}$ W/cm², the underdense plasma is then rather small. For example, LASNEX calculations indicate that the density scale lengths are ~ 10 -30 μ .

Because of the small mass and heat capacity of the underdense plasma, the stimulated reflectivity of high intensity light can rapidly self-limit when account is taken of the momentum and energy deposition into the plasma via the electrostatic waves.¹³ One very simple example of this effect is shown in Fig. 6. This is a plot of back-reflection via the Brillouin instability vs. time as computed with a 1-D hybrid code (fluid electrons and particle ions). In this example, 1.06 μ radiation with an intensity of 10^{16} W/cm² is incident on a 13.5 μ slab of initially uniform plasma with a density $n = 1/3 n_{cr}$ and with an electron temperature of 1 keV. The net amount of plasma is motivated by LASNEX simulations of how much plasma is in the underdense region when the laser pulse reaches high intensity in the experiments with $I \sim 10^{16}$ W/cm². Note that the stimulated reflectivity turns off

Fig. 6

An example of the back-reflection due to the Brillouin instability vs. time as computed with a 1-D code. Plasma conditions and light intensity are given in the text. The boundary conditions model contact with a reservoir of plasma at the initial temperature as detailed in Ref. 13. The time is scaled from an ion-electron mass ratio of 100 to a ratio of 3600 assuming a square root dependence.

in a time of $\lesssim 20$ ps. This is in agreement with estimates of how long it takes for the reflecting light to push the underdense plasma back, modifying its size and gradients enough to turn off the large reflection. Although this model is one-dimensional, clearly a similar behavior is expected for side-scattered light.

As shown in Fig. 7, the back reflection in the experiments with $I \geq 10^{16}$ W/cm² is time-dependent² (at least in most of the shots). Note that a sizeable reflectivity persists for ~ 20 ps, consistent with the simple model. The largest uncertainty in these calculations is just how much relatively cold plasma ($\theta_e \sim 1$ keV) exists in the underdense plasma when the light pulse reaches a high intensity. This depends on both the transport and the form of the heated distributions, both of which are poorly known in detail.

Alternative explanations for a time-dependent reflectivity can be constructed. For example, Brillouin scatter can be small if the entire underdense plasma is at an extremely high temperature. The time-dependence could then be due to a break-up of the

Fig. 7

The time dependence of the reflection back through the focusing lens measured in several experiments with different intensities. The solid line denotes the incident pulse and the dashed line denotes the reflected pulse. These pulses are not plotted on the same intensity scale.

critical surface on a hydrodynamic time scale, reflecting light into a larger solid angle. An improved understanding of the role of induced scattering awaits future experiments with more extensive regions of underdense plasma in which the temperatures and gradients are better characterized.

HEATED ELECTRONS

We will conclude with a few comments on the heated electron energies. The measured X-ray spectra show a hard component. To calculate this spectrum, the LASHEX simulations^{17,18} (and simple estimates) show that a large fraction of the absorbed energy must be put into rather hot electrons. For example, the hot electron energy employed in these calculations are ≈ 10 keV for $I = 10^{15}$ W/cm² and ≈ 40 keV for $I = 10^{16}$ W/cm². The 2-D particle simulations show that such energies are approximately what is expected from the collective absorption mechanisms we have discussed.

An estimate will help to clarify this point. From the cold wave breaking condition, the mean energy of the fast electrons produced by the resonantly-excited waves is $\sim 2^{-4}$

$$\langle \epsilon_e \rangle \sim \frac{1}{2} mc^2 \left(\frac{v_{os}}{c} \right) \sqrt{k_0 L} \phi(\tau)$$

where v_{os} is the quiver velocity of an electron in the free-space value of the laser field and $\phi(\tau)$ is the resonance function previously mentioned. For $\phi(\tau) \sim 1$ and taking $k_0 L \sim 2$ (locally steepened profile), $\langle \epsilon_e \rangle \sim 350 v_{os}/c$ (keV). For $I = 10^{16}$ W/cm², $\langle \epsilon_e \rangle \sim 35$ keV. For $I \leq 10^{16}$ W/cm², $\langle \epsilon_e \rangle \sim 12$ keV. Note in particular that these predictions would be much larger if one neglected the nonlinear profile steepening which is self-consistently included in the particle simulations.

Finally, we emphasize that relating the X-ray spectrum to heated distributions is rather indirect, depending not only on the heated electron spectra but also upon the space and energy dependent transport inhibition. LASNEX calculations show that different energy X-rays dominantly come from quite different regions of the plasma. Hence more detailed comparisons between calculations and experiment depends on an improved understanding of the heat transport (and the various mechanisms for its inhibition). This would enable, for example, the use of improved boundary conditions in the particle simulations. In view of such uncertainties it is premature to deduce details of the shape of the heated distributions from the data.

SUMMARY

In conclusion, we have discussed some recent measurements of the absorption of 1.06 μ radiation with focused intensities in the range $10^{15} - 10^{17}$ W/cm². The magnitude of the measured absorption efficiency is in the range of 30-40%, and the scattered light has an asymmetry with respect to the polarization vector of the incident light. The data was shown to be generally inconsistent with classical inverse bremsstrahlung. The magnitude of the absorption, the polarization dependence of the scattered light, and the heated electron energies were shown to compare favorably with recent calculations of light absorption via plasma wave generation. Improved models necessary for more quantitative comparisons were briefly discussed with reference to details of the data. Although the quantitative understanding of collective plasma effects on light absorption is at an early stage, correlations between the calculations and the data confirm that such processes play an important role.

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