Surface energy transport following relativistic laser-solid interaction

H. Langhoff,1,a) B. T. Bowes,1 M. C. Downer,1,b) Bixue Hou,2 and J. A. Nees2

1Department of Physics, University of Texas at Austin, Austin, Texas 78712, USA
2Center for Ultrafast Optical Science, University of Michigan, Ann Arbor, Michigan 48109, USA

(Received 1 April 2009; accepted 3 June 2009; published online 6 July 2009)

A planar Al target is excited by a 25 fs laser pulse focused to intensity up to $3 \times 10^{18}$ W/cm$^2$ in a $\sim 1$ $\mu$m radius spot; subsequent heat propagation along the target surface, imaged by a delayed probe pulse, appears as a roughly circular area of reduced reflectivity centered on the pump spot, that expands to as much as $12 \pm 3$ $\mu$m in radius within 500 fs. We present a semiempirical model in which the pump laser pulse drives hot electrons into the target via collisionless interactions. A return current heats the target and, above a critical temperature, includes runaway electrons that return to the surface before dissipating their energy. Ultrafast radial expansion of the heated surface layer is explained by lateral diffusive motion of returning runaway electrons oscillating across the target surface layer confined by space charge. Isotropy of the observed expansion is consistent with dominance of resonance absorption over $j \times B$ heating, indicating prepulse heating is important.

© 2009 American Institute of Physics. [DOI: 10.1063/1.3158950]

I. INTRODUCTION

The physics of relativistically intense, ultrashort laser pulses interacting with solid targets is relevant to fast ignition of laser fusion,1,2 to generation of ultrashort pulsed x rays,3,4 and relativistic ion beams,5 and to astrophysics.6 Such laser pulses can deposit as much as half their energy into suprathermal electrons (tens of keV to several MeV),7 which become the dominant carriers of energy out of the initially photoexcited volume. Such fast electrons have been extensively explored in several recent papers by observing them in photoexcited volume. Such fast electrons have been extensively explored in several recent papers by observing them in photoexcited volume. However, strong electric and magnetic fields in the interaction region can also significantly redirect hot electron currents and energy flow along the surface of the target, particularly with obliquely incident laser pulses.2,12,13 Femtosecond-time- and micron-space-resolved measurements of the lateral spread of energy and detailed theoretical description of the transfer of energy from the laser pulse to the surrounding target surface are lacking.

In a previous letter,14 we reported preliminary results of a femtosecond microscopy experiment aimed at providing such a measurement. An aluminum target surface was intensely irradiated by an ultrashort pump pulse focused to peak intensity $I_{pu} \approx 10^{18}$ W/cm$^2$ within a spot radius $w_0 \sim 1$ $\mu$m. Time-resolved images of the surface obtained with an ultrashort probe pulse showed sharply reduced reflectivity within the directly heated area. Moreover, as $I_{pu}$ approached relativistic values, the radius of the dark hot region expanded threefold within $\sim 100$ fs. By way of qualitative explanation, we pointed out that for $I_{pu} \approx 10^{18}$ W/cm$^2$, the mean free paths for both collisional and radiative transport were expected to become comparable to $w_0$. Thus a rapid burst of ballistic transport outward from the photoexcited area could be expected. However, no quantitative model of the lateral heat transport physics was presented.

The primary purpose of this paper is to present a theoretical model of the transport of energy from the pump laser laterally to the surrounding target surface in order to explain the experimental results. We adopt a semiempirical approach to render the underlying physics as transparent as possible. Each step of the transport process—production of suprathermal electrons, generation of a return current, and interaction of runaway return electrons with the target surface—is calculated separately, with clear assumptions and approximations at each step. Nevertheless, intermediate calculated results (e.g., hot electron temperature, self-generated electrostatic fields, etc.) are benchmarked against results of independent particle-in-cell (PIC) simulations where available. This model satisfactorily explains previous experimental results14 and predicts that the scale length of lateral heat transport should increase nonlinearly at higher pump intensities in the same focus geometry. To test this prediction, we upgraded the $\lambda^3$ laser system to produce focused intensity $I_{pu} \approx 3 \times 10^{18}$ W/cm$^2$ in the same spot size on the Al target surface. A second purpose of this paper is to present results from this new round of experiments. Consistent with our model, we obtain clear evidence that the “dark” hot region can grow to a radius as much as four times larger than observed at $I_{pu} \approx 1 \times 10^{18}$ W/cm$^2$ on a similar time scale. Also consistent with our model, we find that the size of the imaged hot region is much more sensitive to small shot-to-shot intensity fluctuations, arising either from the laser system itself or wobble of the rotating target plane, than at lower pump intensities. At $\Delta t \approx 200$ fs, hot regions with radii as large as 10 $\mu$m were observed on selected shots. A third purpose of this paper is to describe the experiment in more detail than was possible in a brief letter.

---

---
The remainder of Fig. 1 summarizes, for completeness, previously reported results. The images at the bottom of Fig. 1(a) show examples of reflected probe profiles before arrival of the pump (left), immediately after arrival of the pump (center) showing an area of reduced reflectivity equal in size to the pump spot, and at Δt > 1 ps (right), where the dark heated region was observed to expand as \( I_{pu} \) approached relativistic values. Figure 1(b) shows that for subrelativistic pump intensity (\( I_{pu} \approx 3 \times 10^{17} \) W/cm\(^2\)), the dark heated region remained within the directly photoexcited pump spot for all Δt investigated (right). For weakly relativistic pump intensity (\( I_{pu} \approx 1 \times 10^{18} \) W/cm\(^2\)), on the other hand, the heated region grew to approximately three times the diameter of the pump spot within 500 fs (left).

The data points in Fig. 2 plot the measured diameter of the hot dark region versus Δt for \( I_{pu} = 10^{18} \) W/cm\(^2\). Error bars reflect variations in diameter measured across different cross sections as well as fluctuations from run to run. The calculated curve is discussed in Sec. III. The final expanded area was close to the source area for keV x rays generated under the same conditions, observed by time-integrated x-ray shadowgraphy.

The λ\(^3\) laser was subsequently upgraded with an additional amplifier that produced approximately five times higher energy per pulse focussable to \( 2w_0 = 2.0 \) μm on the target surface using adaptive optics, resulting in approximately three times higher peak intensity. Third-order autocorrelation measurements revealed no prepulses higher than \( \sim 10^{-3}I_{pu} \) within \( \sim 1 \) ns of the main pulse. The results of femtosecond microscopy measurements differed in three ways from the earlier results, as illustrated in Fig. 3. First, surface second-harmonic generation (SHG) by the \( p \)-polarized pump pulse became stronger. On most shots, a small portion of this 400 nm light scattered nonspecularly...
No changes in probe reflectivity were observed before $\Delta t = 0$, indicating that the main pump pulse, rather than prepulses, was primarily responsible for changes at $\Delta t \geq 0$. Nevertheless, this does not rule out the possibility that prepulses preheated the surface. Third, however, we observed that the size of the imaged dark region fluctuated three to four times more from shot to shot than the $\sim 20\%$ fluctuations observed at lower pump intensities. The main source of these fluctuations was shot-to-shot fluctuations in laser pulse energy. Wobble of the target plane as the sample rotated, which had to be carefully limited to less than the few microns depth of focus of the pump pulse using precision micrometer adjustments on the sample mount, may have contributed additional periodic variations in pump intensity. Since pump energy fluctuations and sample plane adjustment procedures were identical in both the low- and high-intensity measurements, we conclude that the energy transport physics itself became more sensitive to fluctuations in absorbed pump intensity at higher $I_{pu}$. The increased amount and lateral range of debris from foregoing laser shots at higher $I_{pu}$ and fluctuations in prepulse heating may have also contributed to fluctuations in the absorbed fraction of $I_{pu}$. Because of these strong fluctuations, we were unable to measure temporal growth of the hot region reliably as in Figs. 1(b) and 2. Thus in Figs. 3(b) and 3(c) we present only an example of one of the largest hot, dark region images observed, corresponding presumably to the largest absorbed $I_{pu}$. The irregular shape of the expanded dark region highlighted by the dashed outline curve in Fig. 3(b) was typical and changed from shot to shot. However, no consistent evidence of anisotropic expansion with respect to the projection of the pump incident plane on the sample surface, shown by a dashed arrow in Fig. 3(b), was observed. Recent time-integrated x-ray shadowgraph measurements with the upgraded $\lambda^3$ laser$^{18}$ revealed x-ray source size several times larger than observed at lower intensity,$^{15}$ qualitatively consistent with the present results.

Negligible pump-induced changes in reflectivity were observed with an $s$-polarized probe. The selective decrease in reflectivity $\Delta R_{p}^{p}$ for the $p$-polarized probe is caused by a combination of resonance absorption (RA)$^{19,20}$ and reduced Fresnel reflectivity in the hot plasma. RA is only efficient after an electron density gradient develops over a scale length $L \approx \lambda_{pr}/2\pi$, which requires expansion of aluminum ions (mass $m_i$ and charge state $Z$) at velocity $v_i = (ZKTe/m_i)^{1/2}$ (Ref. 21) into the vacuum. Indeed we showed in Ref. 14 that the reflectivity drop $\Delta R_{p}^{p}$ near the center of the pump spot ($r < w_0$) was delayed by $\Delta t \approx 100$ fs, consistent with $Z \sim 6$ and $KT_e \sim 1$ keV for $I_{pu} \sim 10^{18}$ W/cm$^2$. Faster response is expected at higher $I_{pu}$. In addition, Fresnel reflectivity can change instantaneously in response to plasma heating for $L \approx 0$, depending on local electron temperature. For $r < w_0$ at $I_{pu} \sim 10^{18}$ W/cm$^2$, keV electron temperatures and universal plasma mirror$^{22}$ properties are expected at $\Delta t = 0$. Under these conditions, Price et al.$^{22}$ observed that the reflectivity of 400 nm light incident normally on Al increased from 80% to 90%, although Bowes et al.$^{14}$ observed no change in the reflectivity of 400 nm probe pulses incident into the probe imaging optics, producing a bright and useful fiducial marker of the location and size of the pump focus near the center of the image on each shot, as evident in all three panels of Fig. 3. Second, we observed on many shots that the dark, hot region expanded dramatically to as much as $12 \pm 3$ $\mu$m radius for $\Delta t \approx 300$ fs, as illustrated in Fig. 3(b).
at $\theta_{pr}=70^\circ$ at $\Delta t=0$. Thus the slightly delayed decrease $\Delta R_p^\parallel$ observed at $r<w_0$ can be attributed mainly to RA. On the other hand, lateral energy transport heats surrounding target material ($r>2w_0$) to temperatures of only tens to hundreds of eV. For these temperatures, Price et al. observed the reflectivity of 400 nm light incident normally on Al decreased to <70%, while several other investigators reported even steeper, instantaneous drops in reflectivity of obliquely incident, $p$-polarized light, prior to onset of delayed RA, from Al targets heated to temperatures in this range. These reflectivity drops result from increased collisional absorption. Changes in $x$-polarized reflectivity were negligible in comparison, consistent with the present findings. Thus grazing incidence, $p$-polarized 400 nm probe light is a sensitive probe of the outer boundary of the heated region that responds without time delay to the deposition of heat.

From analysis of prior results, we estimate our probe to be sensitive to surface temperatures exceeding $kT_{thresh} \sim 40$ eV at the outer boundary. It is thus sensitive to much lower temperatures than x-ray diagnostics. The optical response of other target materials to sub-keV heating varies widely, so the femtosecond microscopy technique employed here would have to be recalibrated based on independent pump-probe measurements of those materials.

### III. SEMIEMPIRICAL MODEL OF LATERAL HEAT TRANSPORT

We model heat transport out of the initially photoexcited volume in three steps: (i) collisionless laser-plasma interactions drive a current of suprathermal electrons into the target, (ii) the space charge field created by these primary electrons drives a return current of secondary electrons that Ohmically heats the target, and (iii) some runaway electrons return to the target surface without colliding, then oscillate and spread diffusively along the target surface. We consider each step in turn.

#### A. Suprathermal electrons

While at low intensities laser absorption in solid targets is mainly collisional, relativistically intense laser pulses couple to plasma electrons predominantly via collisionless mechanisms including RA, vacuum heating (VH), and $j \times B$ heating, which can produce suprathermal (hot) electrons with efficiency approaching 50%. RA and VH accelerate hot electrons primarily normal to the target surface. The $j \times B$ force, on the other hand, accelerates electrons in the direction of the incident pump pulse and thus should heat the surface anisotropically when the pump is obliquely incident. Indeed, the observation of fast electrons propagating along the surface of a target excited at large $\theta_{inc}$ by a relativistically intense pulse was attributed to the action of the $j \times B$ force. On the other hand, the intensity of the surface electron beam diminished rapidly with increasing prepulse intensity. Since we observed no anisotropy in the probe absorption pattern in the present experiment, we presume that RA and/or VH are the dominant mechanisms for production of suprathermal electrons under our conditions. This in turn indicates that prepulses, despite producing no observable probe reflectivity change prior to $\Delta t=0$, may have preheated the surface sufficiently that RA and VH dominate hot electron production.

Although there is no intrinsic reason for hot electrons produced by collisionless processes to possess a well-defined temperature, approximately Maxwellian energy distributions have been widely observed both in PIC simulations and in experiments. Accordingly, in our numerical model we launched hot electrons into the target normal to its surface ($z$ coordinate) with a temperature given by the empirical rule established by Beg, where $kT_{e}^{hot}$ is evaluated in keV, absorbed pump intensity $I_{pu}^{(18)}$ (averaged over the focal spot) in units of $10^{18}$ W/cm$^2$, and laser wavelength $\lambda_l$ in units of microns. A similar relation may be found by analyzing the x-ray data reported by Theobald et al. For numerical calculations, absorbed intensity was taken to be the spatially averaged incident intensity $I_{pu}$ multiplied by measured pump absorption fraction $a=0.5$ and by $\cos \theta_{pu}=1/\sqrt{2}$. Ions were held fixed for the investigated time interval.

Hot electrons propagating into the target are slowed partly by the space charge electric field resulting from the imperfect balance of primary and return current densities and partly by collisions. In the present situation, deceleration by the space charge field dominates. The range of a single primary electron with energy $E=kT_{e}^{hot}$ in eV is where $\sigma_e$ is the electrical conductivity of the target and $e$ is the elementary charge. The $z$ component of the space charge electric field then becomes

$$X(0) \approx \frac{E}{e\tau_{eff}} = \frac{e a I_{pu} \cos \theta_{pu}}{\sigma_e kT_{e,0}^{hot}}$$

(3)

near the surface (0 $< z < \tau_{eff}$). Table I lists calculated values of $\tau_{eff}$, $X(0)$, and $\sigma_e$ from the model for selected incident pump intensities.

In the actual modeling, the Maxwellian energy distribution of the primary electrons was taken into account. The angular spread of the hot electrons was approximated by as-

<table>
<thead>
<tr>
<th>$I_{pu}^{(18)}$ (1 × 10$^{18}$ W/cm$^2$)</th>
<th>$\tau_{eff}$ (µm)</th>
<th>$X(0)$ (GV/m)</th>
<th>$\sigma_e(0)$ ($\times 10^4$Ω m$^{-1}$)</th>
<th>$\sigma_e(z)$ ($\times 10^4$Ω m$^{-1}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.5</td>
<td>1.51</td>
<td>24.0</td>
<td>4.8</td>
<td>5.1</td>
</tr>
<tr>
<td>1</td>
<td>1.83</td>
<td>23.3</td>
<td>7.8</td>
<td>6.0</td>
</tr>
<tr>
<td>2</td>
<td>2.44</td>
<td>24.1</td>
<td>12.1</td>
<td>8.2</td>
</tr>
<tr>
<td>4</td>
<td>3.47</td>
<td>23.5</td>
<td>19.9</td>
<td>11.7</td>
</tr>
<tr>
<td>8</td>
<td>4.35</td>
<td>22.1</td>
<td>33.7</td>
<td>14.6</td>
</tr>
<tr>
<td>16</td>
<td>5.82</td>
<td>20.5</td>
<td>57.9</td>
<td>20.9</td>
</tr>
</tbody>
</table>

For single electron with $kT_{e}^{hot}$.

Author complimentary copy. Redistribution subject to AIP license or copyright, see http://php.aip.org/php/copyright.jsp
assuming the radius of the electron beam expanded like the radius \( w(z) = w_0(1 + z^2/z_0^2)^{1/2} \) of the pump beam propagating in vacuum, where \( z_0 = \pi w_0^2/\lambda_{pu} \) is its Rayleigh length. The space charge electric field is then given by

\[
X(z) = \frac{X_0}{[1 + eX_0z_0 \tan^{-1}(z/z_0)/kT_e^\text{bot}][1 + z^2/z_0^2]},
\]

(4)

**B. Return current**

The return current is responsible for heating the target. Ion temperature \( T_i \) over the heated zone of diameter \( 2w(z) \) may be calculated from

\[
d(c_i/\theta) = \sigma_Z^2(z,T_e^\text{bot}).
\]

(5)

For the integration the temperature dependence of the heat capacity \( c_i(\theta) \) and conductivity \( \sigma_Z(\theta) \) were taken into account. Expressions for \( c_i(\theta) \) were obtained by solving a set of Saha equations describing the different ionization states of aluminum. These calculations are consistent with the empirical expression \( Z = 13(\theta_{eV})/(125 + \theta_{eV}) \) for the average charge state \( Z \) at aluminum solid density, which yields estimated \( Z \approx 7 \) at our highest experimental intensity. For \( \sigma_Z(\theta) \), experimental values reported by Milchberg et al.\(^{23} \) were used for \( \theta \approx 70 \text{ eV} \) and Spitzer’s formula for conductivity\(^{32} \) at \( \theta > 70 \text{ eV} \). The values agreed with the calculations by Lee and More.\(^{33} \) The average energy that electrons of the return current gain in the electric field \( X(z) \) is determined by their mean free path,

\[
\Delta(\theta,T_e) = \frac{1}{n_Z(\theta)\sigma_{\text{scatt}}(T_e)},
\]

(6)

where \( n_Z(\theta) \) denotes the density of aluminum ions with average charge \( Z \) at temperature \( \theta \) and \( \sigma_{\text{scatt}}(T_e) \) is the cross section for momentum transfer from returning electrons of energy to ions. For small 4\( T_e \), \( \sigma_{\text{scatt}} \) was computed by Mayol and Savat.\(^{34} \) For larger 4\( T_e \), the electron-ion cross section \( \sigma_{ci} \) for elastic scattering compiled by Anders\(^{35} \) is a more appropriate approximation. The energy loss of the electrons by inelastic collisions was calculated using semiempirical expressions given by Voronov\(^{36} \) for different ionization states \( Z(\theta) \) of aluminum ions in the plasma. The ionization rates decrease rapidly with rising \( Z(\theta) \) and with the kinetic energy of the electrons.

Figure 4(a) shows results of the calculations of \( \theta \) at the surface at the end of the laser pulse versus pump intensity \( I_{\text{pu}}^{\text{(18)}} \). Plasma electrons, which equilibrate rapidly with ion temperature \( \theta \), reach the threshold value (\( \sim 40 \text{ eV} \)) detectable by probe reflectivity at \( I_{\text{pu}}^{\text{(18)}} \sim 0.3 \), consistent with the results shown on the right-hand side of Fig. 1(b). At higher intensity, temperature rises nonlinearly because of the complex dependence of \( kT_e^\text{bot} \), \( \varepsilon_{\text{eff}} \), \( \sigma_c \), and \( X \) on intensity. The results are consistent with similar calculations by Price et al.\(^{22} \) using the LASNEX code. Table II shows the calculated depth profiles \( \theta(z) \) and \( X(z) \) of temperature and electric field, respectively, for selected intensity \( I_{\text{pu}}^{\text{(18)}} = 1 \).

**C. Runaway electrons**

With increasing intensity \( I_{\text{pu}}^{\text{(18)}} \) and time \( \theta \) the temperature reaches a critical value \( \theta_{cr} \) at which the momentum normal to the surface that a returning electron gains in the electric field \( X(z) \) within mean free path \( \Delta \) exceeds the momentum exchanged by collisions, causing a beam of runaway electrons to develop.\(^{32,37} \) As discussed above, both elastic and inelastic collisions can influence \( \Delta \) and the momentum exchanged, depending on plasma conditions. In either case, as temperature increases above \( \theta_{cr} \), collision probability rapidly decreases while \( \Delta \) and thus momentum gain within \( \Delta \) increase, resulting in continuous acceleration toward the surface, the primary feature of runaway electrons.\(^{32,37} \) For the present

---

**TABLE II.** Calculated depth \( (z) \) profiles of temperature \( [\theta(z)] \) and space charge electric field \( [X(z)] \) for incident pump intensity \( I_{\text{pu}}^{\text{(18)}} = 1 \).

<table>
<thead>
<tr>
<th>( z ) ((\mu\text{m}))</th>
<th>( \theta(z) ) ((\text{eV}))</th>
<th>( X(z) ) ((\text{GV/m}))</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.25</td>
<td>84.6</td>
<td>23.3</td>
</tr>
<tr>
<td>0.75</td>
<td>71.7</td>
<td>23.3</td>
</tr>
<tr>
<td>1.25</td>
<td>54.8</td>
<td>23.1</td>
</tr>
<tr>
<td>1.75</td>
<td>38.3</td>
<td>17.8</td>
</tr>
<tr>
<td>2.25</td>
<td>25.8</td>
<td>12.4</td>
</tr>
<tr>
<td>2.75</td>
<td>16.9</td>
<td>8.5</td>
</tr>
<tr>
<td>3.25</td>
<td>10.9</td>
<td>5.7</td>
</tr>
</tbody>
</table>

---
diffusive motion originates from electron-ion scattering during anisotropic surface energy transport above a critical electron energy loss rates in aluminum given by Voronov,36 noting that runaway electrons spend only half their time in aluminum. The electron energy loss heats the surface ions, which equilibrate rapidly with surface plasma electrons, causing onset of substantial probe absorption when temperature exceeds \( kT_{\text{thresh}} \sim 40 \) eV, as discussed in Sec. II.

Figure 4 shows the calculated diameter \( 2r_{\text{diff}} \) for \( t > \tau \) as a function of \( \mu_{\text{pu}}^{(18)} \). The behavior of \( 2r_{\text{diff}} \)(\( \mu_{\text{pu}}^{(18)} \)) indeed mirrors that of runaway electron temperature \( kT_{\text{eff}}(\mu_{\text{pu}}^{(18)}) \) shown in Fig. 4(b); expansion beyond the initially photoexcited area begins to occur at \( \mu_{\text{pu}}^{(18)} \sim 0.5 \), then increases sharply for \( \mu_{\text{pu}}^{(18)} > 1.0 \). Experimental values of \( 2r_{\text{diff}} \) are plotted for comparison. In view of the uncertainties in the data and in the assumed parameters for the calculations, reasonable agreement is obtained. As important as individual numerical values of \( 2r_{\text{diff}} \) is the overall trend in the slope of \( 2r_{\text{diff}} \)(\( \mu_{\text{pu}}^{(18)} \)). The low slope for \( \mu_{\text{pu}}^{(18)} < 1.0 \) is consistent with observation of stable, reproducible values of \( 2r_{\text{diff}} \)(\( \mu_{\text{pu}}^{(18)} \)) in that range in Ref. 14, whereas the high slope for \( \mu_{\text{pu}}^{(18)} > 1.0 \) is consistent with the observation of large shot-to-shot fluctuations in \( 2r_{\text{diff}} \). For example, from Fig. 5 fluctuations of \( \pm 50\% \) around the value \( I_{\text{pu}}^{(18)} = 2.0 \) would lead to fluctuations over the range 4 \( < 2r_{\text{diff}} < 30 \) \( \mu \text{m} \) in the observed diameter of the hot dark region.

In the range 0.5 \( < \mu_{\text{pu}}^{(18)} < 1.0 \) the time evolution \( r_{\text{diff}}(t) \) from Eq. (8) can be compared with experimental data from Ref. 14. Figure 2 shows this comparison for \( I_{\text{pu}}^{(18)} = 1.0 \). Very good agreement is obtained. The close agreement shown may be somewhat fortuitous, however, in view of uncertainties in both experiment and calculation. Nevertheless, the approximate time scale and the total range of expansion are correctly modeled. The latter are the significant points of agreement to be expected for this level of computation.

**IV. DISCUSSION**

Recent literature on solid targets irradiated at relativistic intensity at oblique incidence has strongly emphasized calculation12 and measurement2,13 of anisotropic hot electron transport out of the initially photoexcited volume. Indeed this anisotropy was a key element of fast ignition schemes employing hollow gold cones to channel laser-generated surface hot electrons into a fusion target.7 The present observation of approximately isotropic radial transport appears to run counter to the expectations of this previous literature, but at the same time provides a framework for understanding how isotropic transport can occur even with oblique incidence, relativistically intense excitation. In fact the isotropy observed here may be entirely consistent with previous literature. In the theory of Nakamura et al.,15 the most strongly anisotropic surface energy transport occurs above a critical pump incidence angle (\( \sim 70^\circ \)) larger than used in the present experiment.

![Diagram of surface region with temperature exceeding \( kT_{\text{thresh}} = 40 \) eV vs pump intensity. Data points: measured diameter of region of reduced probe reflectivity. Solid curve: Estimated diameter after stoppage of runaway electrons from model calculations. Upper left inset: drawing of diffusive motion of runaway electrons.](image_url)
experiments. At smaller angles energy penetrates into the volume as in the picture proposed here. Similarly, in the experiments of Li et al., fast surface electrons are observed prominently only for $\theta_{pu} \sim 70^\circ$. A merit of reflectivity experiments is that electrostatic forces that can influence the trajectory of electrons escaping the target play no role.

Several improvements and extensions of the femtosecond microscopy method presented here may help link these disparate observations. First, pump pulses of higher peak-to-background contrast, higher energy, and looser focus should be employed. Higher contrast will suppress target preheating further and increase the importance of $j \times B$ heating, and thus of anisotropic transport. Indeed multiterawatt laser systems with contrast exceeding $10^{11}$ have recently been developed. Pump pulses of higher energy but looser focus will help mitigate sensitivity to target plane wobble. Taken together, these improvements should enable explicitly femtosecond-time-resolved measurement of radial transport following ultrarelativistic irradiation. Second, femtosecond microscopy measurements on the front surface of the target should be correlated with direct characterization of hot electrons from the back surface. For example, Cho et al. recently observed two distinct lobes of coherent transition radiation (CTR) from the back surface of a thin Al foil irradiated intensively at oblique incidence on the front surface. One lobe originated from hot electrons propagating normal to the target surface, generated by VH and RA, the other from hot electrons co-propagating with the obliquely incident pump laser pulse, produced by $j \times B$ heating. Comparison of the intensity of the two CTR lobes enables immediate evaluation of the relative importance of $j \times B$ heating. To implement this extension the conflicting needs for a thick target to facilitate maintenance of a stable, optically flat target plane and a thin ($\sim 10 \mu m$) target for back surface CTR must be reconciled. Use of pump pulses with larger depth of focus than in the present experiments will thus be critical in this regard. Third, femtosecond microscopy should also be correlated with time-integrated x-ray and electron emission measurements using the same target and conditions. The former provides a complementary measurement of hot electron transport, but should yield lateral x-ray source sizes smaller than reduced reflectivity regions, since reflectivity is sensitive to temperatures down to $kT_{thres} \sim 40 \text{ eV}$. The latter can potentially measure $kT_{el}^{hot}$ independently and help to refine the empirical intensity scaling law [Eq. (1)] used in modeling. Finally, with the previous improvements in place, the angle of incidence $\theta_{pu}$ should be systematically varied. Previous theory and experiment both suggest that anisotropic production of surface fast electrons increases nonlinearly with $\theta_{pu}$ and can become the dominant transport process above a certain critical angle. With systematically varied $\theta_{pu}$ it should be possible to document the transition from isotropic to anisotropic hot electron transport explicitly.

The above improvements combined with calibration of $\Delta R_p^b$ based on independent measurements of the optical reflectance properties of intensely irradiated targets could potentially enable temperature distributions $\theta(r, \Delta t)$ around an intense laser focus to be measured, in addition to tracking the periphery of the hot region as in this work. The high space- and time-resolution of femtosecond microscopy measurements should also make them a valuable benchmark of large-scale PIC simulations of intense laser-solid interactions that complements measurements of x-rays, electrons, and ions. In particular, the analysis of Sec. III suggests that the radius of the heated region is a sensitive and direct indicator of the temperature and number of runaway electrons, a quantity only indirectly related to these other measurements.

V. CONCLUSIONS

We presented femtosecond microscopy measurements of the reflectivity of a planar Al surface within and surrounding a spot of radius $w_0 \sim 1 \mu m$ in which a relativistically intense, obliquely incident pump pulse was absorbed. For $I_{pu} > 10^{18} \text{ W/cm}^2$ reflectivity is sharply reduced within several hundred femtoseconds out to a radius as much as $10w_0$. We explain these results using a semiempirical model consisting of three steps: (1) production of suprathermal electrons by VH and/or RA, (2) production of a return current driven by the space charge field created by the primary suprathermal electron current, and (3) lateral diffusive motion of the runaway electron component of the return current. Thus the absorbed energy of the pump laser is partitioned into a hot ($0.1 < \theta < 0.5 \text{ keV}$) central region with the diameter $2w_0$ of the pump focus and depth of order $1.5 < z_{eff} < 3 \mu m$ created by the collisional component of the return current and a warm ($\theta \sim 50 \text{ eV}$) sheath of several microns radius confined to the surface, created by the runaway electron component of the return current that diffuses rapidly along the surface, depositing energy by inelastic collisions. Improved and expanded versions of femtosecond microscopy experiments should provide a key benchmark for codes used to model x-ray and ion generation, laboratory astrophysical phenomena in intensely irradiated solid targets.

ACKNOWLEDGMENTS

This work was supported by the National Science Foundation Grant No. PHY-0114336, U.S. Department of Energy Grant No. DE-FG03-96ER40954 and Army Research Office Grant No. W911NF-07-1-0056.


25. L. Spitzer, Jr., Physics of Fully Ionized Gases (Interscience, New York, 1962), Sec. 5.4.


