Comparative Reflectance Measurements on Laser-Produced Plasmas at 1.06 and 0.53 μm

A. G. M. Maaswinkel, K. Eidmann, and R. Sigel
Projektgruppe für Laserforschung der Max-Planck-Gesellschaft zur Förderung der Wissenschaften e.V., D-8046 Garching bei München, Germany
(Received 28 March 1979)

The total and specular reflectances of 1.06- and 0.53-μm laser pulses were measured on planar copper targets. The measurements were made with 30-ps pulses and an intensity on target of 10^{14} W/cm^2. The direction of polarization and angle of incidence (0°–70°) were varied. Only a weak dependence of reflectance on wavelength is found; at both wavelengths reflectance is governed by the characteristic angle and polarization dependence of resonance absorption.

Absorption and reflection of intense laser radiation on solid targets is a basic problem in laser-fusion studies. A crucial parameter in this respect is the wavelength of the laser. Up to now, most laser plasma experiments have been made with the CO_2 (10.6 μm), Nd (1.06 μm) and iodine lasers (1.315 μm). No systematic studies, however, have been undertaken in the shorter-wavelength region, although considerable effort is being made to develop visible and near-uv lasers. In this Letter we present measurements of the total and specular reflectances at 1.06 and 0.53 μm and pulse duration of 30 ps. Since recent measurements[1-4] suggest that resonance absorption is an important absorption mechanism, we have studied in detail the dependence of reflectance on the angle of incidence and polarization of the incident laser light. Much attention has been paid to ensuring identical experimental conditions at both wavelengths to allow comparison of the results.

The laser is a Nd:YAlG (neodymium-doped yttrium aluminum garnet) system delivering 0.7 J in 30 ps. A high-contrast Pockelscell-Glan prism combination (extinction ratio 2×10^5) guarantees a clean pulse (prepulses are < 6 μJ). For measurements at 1.06-μm isolation against back-reflected light and rotation of the polarization is achieved by two Faraday rotators; the resulting energy on target is 300 mJ. For 0.53-μm generation we use a potassium dihydrogen phosphate type-II crystal of 51 mm in diameter and 17 mm in length. Its conversion efficiency is 40%, giving 200–300 mJ on target. The polarization here is rotated with a $\lambda/2$ plate. At both wavelengths the polarization is better than 95%. The laser light is focused on target by an achromatic 1:8/400-mm lens. The focal spot was measured by focusing through pinholes.[5] It contains 50% of the energy in a spot 40 μm in diameter, resulting in an averaged intensity on target of 3×10^{14} W/cm^2. Our 30-ps pulse, the total angle of the lens of 7°, and a Rayleigh range of 1 cm will enable us to approach a plane geometry in our experiments. Note that with $t=30$ ps and an estimated temperature of $T_e=400$ eV the plasma expands over a distance $L_H=t(Zk_BT_e/m_i)^{1/2}\approx 4$ μm during the laser pulse, i.e., a distance much smaller than the focal-spot diameter. The experiments were done on Cu targets, made from highly polished hemicylindrical steel rods with 2-μm-thick Cu coatings.

It is known that one obtains in general, upon reflection from a plane target, three components of reflected laser light with characteristic angular distributions[6]: (i) a well-collimated, specular beam (R_{spec}) with a divergence similar to that of the incident beam, (ii) diffusely scattered laser light (R_{diff}) with a very broad angular distribution, and (iii) so-called collimated backscatter through the focusing lens due to Brillouin backscattering. In our case the last component was found to be negligible ($<10^{-2}$), at least for angles of incidence larger than 3.5° (the half-angle of the lens) where it can be discriminated against the specular beam. This is probably due to our short pulse duration and correspondingly small plasma size where no substantial amplification of the backscattered component can occur. With an Ulbricht spherical photometer as described previously in[4] we measured first the total reflection coefficient $R_{tot}{}^{s,p}=R_{spec}{}^{s,p}+R_{diff}$; s,p stands for the polarization of the beam. Here we integrate over a solid angle of 4π steradians; the light reflected through the lens and the incident energy were measured separately with photodiodes. We improved the accuracy of the sphere by using four Si p-i-n diodes as detectors; their average signal is accurate to $<5\%$. The response of the sphere is independent of the incident energy from 0 to 300 mJ.[5] Its calibration was checked before and after each series of measurements.

In order to know the contribution of the specular and diffuse components, we have to measure one of them separately. We measured $R_{spec}{}^{s,p}$

with the Ulbricht sphere removed. Tests with burnpaper showed immediately that a well-defined specular beam is produced upon reflection. Quantitatively its energy was measured with a Gentec detector. The solid angle covered by this detector. The solid angle covered by this detector was varied using a diaphragm of variable diameter in the range 2×10^{-2} to 1.2×10^{-1} sr (14°–38° one angle). This did not change the detector signal. Hence the energy of the specular beam is indeed contained in a divergence angle similar to that of the incident beam; on the other hand, we can conclude that the diffusely scattered light has a much broader angular distribution. This measurement of the specular beam energy together with the Ulbricht-sphere measurements then allows us to calculate $R_{\mathrm{diff}}{}^{s,p} = R_{\mathrm{tot}}{}^{s,p} - R_{\mathrm{spec}}{}^{s,p}$.

The result of the reflectance measurements are plotted in Figs. 1 and 2. The data points are mean values for at least ten shots up to twenty shots. The measurements were made with the focus on target, with intensity $(1-3)\times10^{14}$ W/cm². At 1.06 μm (Fig. 1) a maximum absorption of 0.43 occurs for p-polarized light at an angle of incidence of $\vartheta = 20°$. There is a good agreement between this curve and that obtained by Godwin *et al.* in Ref. 4. We notice that 30% of the energy is still absorbed at $\vartheta = 0°$.

At $\lambda = 0.53$ μm, absorption is always somewhat higher. At $\vartheta = 0°$ it is about 0.40, similar to that in Ref. 7; it reaches a maximum of 0.62 for p polarization also at 20° incidence. At both wavelengths a pronounced difference between s- and p-polarized radiation is evident. It is also clear from Figs. 1 and 2 that in general a strong component of diffusely scattered light is present. It dominates reflection losses at normal incidence (0°), but decreases with increasing angle of incidence. Its intensity depends only weakly on the polarization of the incident beam. A comparison of Figs. 1 and 2 shows that the diffuse component at a fixed angle (22°) is reduced for $\lambda = 0.53$ μm as compared to 1.06 μm (at $\lambda = 1.06$ μm and $\vartheta = 0°$ we have $R_{\mathrm{diff}}{}^{s,p} = 0.55$). Additional measurements have shown that R_{diff} decreases with decreasing intensity.

The observed angle and polarization dependence is believed to be due to resonance absorption.[2,4] A complication arises from the presence of diffuse scattering which is not taken into account in the usual theory of resonance absorption. In this situation a fractional absorption due to resonance absorption may be defined as $f_{\mathrm{res}} = (R_{\mathrm{spec}}{}^{s} - R_{\mathrm{spec}}{}^{p})/R_{\mathrm{spec}}{}^{s}$. With this definition only the fractional energy absorbed out of the *specular* beam due to the polarization effect is attributed to resonance absorption. In Fig. 3 the data of Fig. 2 are plotted in this way and compared with a theoretical curve as calculated by Ginzburg.[8,9] There is fair agreement up to the maximum $f_{\mathrm{res}} \approx 0.5$ for $\vartheta = 20°$, but for larger angles the experimental absorption is larger than calculated. The same observation

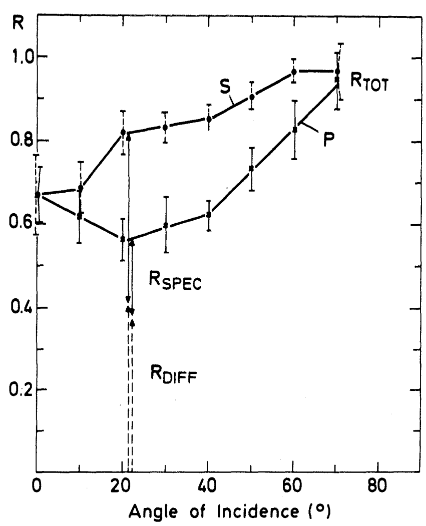

FIG. 1. Reflectance R_{tot} into 4π steradians as a function of the incidence angle for s- and p-polarized light and $\lambda = 1.06$ μm. Specular reflectance R_{spec} measured at $\vartheta = 22°$. Intensity on target is $(1-3)\times10^{14}$ W/cm².

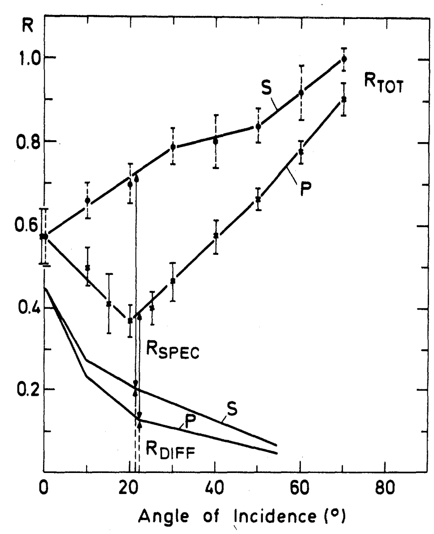

FIG. 2. Reflectance R_{tot}, and R_{spec} at $\vartheta = 0°$, 10°, 22°, and 55° for $\lambda = 0.53$ μm. Intensity is $(1-3)\times10^{14}$ W/cm².

was made more indirectly by Balmer *et al.*[3] From Fig. 1 one obtains $f_{\rm res} = 0.52$ at 22° in agreement with the result at $\lambda = 0.53$ μm. We have also made another set of measurements for $\lambda = 1.06$ and 0.53 μm at 10^{11} W/cm^2. The characteristic polarization dependence remains unchanged, though reflectance decreases, possibly as a result of collisional absorption (for example maximum total absorption is 0.65 at $\lambda = 1.06$ μm and 0.82 at $\lambda = 0.53$ μm). If all these results are plotted as in Fig. 3 they coincide within the error bars with the experimental curve shown there. Thus we have the important result that the resonant part of absorption is *independent of wavelength and intensity* as it is expected to be. Maximum absorption occurs for both wavelengths at $\vartheta_m = 20°$. One may obtain[8] from this angle the scale length L of the density gradient by $(2\pi L/\lambda_0)^{1/3} \sin\vartheta_m \approx 0.7$. One finds $L \approx \lambda_0$ *both* wavelengths, in agreement with previous observations at $\lambda = 1.06$ μm.[1-4]

Resonance absorption in a smooth plasma layer cannot explain the observed absorption at normal incidence. If total absorption at 1.06 μm would be due to collisional absorption, it should increase more strongly than observed for 0.53 μm. Estimates with[10,11] $T_e = 400$ eV and $L = \lambda$ show that an absorption of 10% at 1.06 μm and 20% at 0.53 μm may be attributed to collisional absorption; the major source of absorption most likely has other causes. The observation of strong diffuse scattering at normal incidence suggests that absorption is connected with the underlying "roughness" of the reflecting plasma layer. As has been discussed in Thomson, Kruer, and Langdom[12] resonance absorption may account for nearly all of the absorption in the presence of roughness.

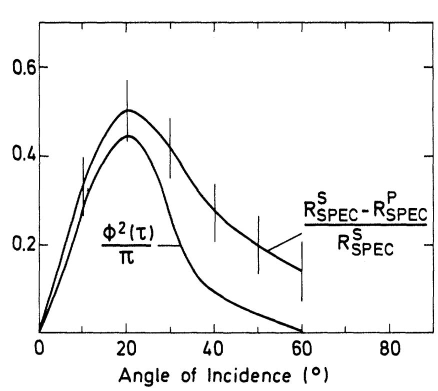

FIG. 3. Comparison at $\lambda = 0.53$ μm between theoretical resonance absorption $\varphi^2(\tau)/\pi$ and the experimental curve obtained from Fig. 2.

Also it seems important to remember that coupling to surface plasmons leads to enhanced light absorption on rough metal surfaces.[13] This effect which is bascially similar to resonance absorption is of interest at least as a limiting effect for a plasma with a very steep density gradient. In contrast to a metal surface, roughness in a laser-produced plasma is, however, generated by the interaction itself and therefore the situation is extremely complicated. At the present time, the cause and importance of diffuse scattering in the context of resonance absorption is difficult to assess; it still surprises us that the predicted behavior of resonance absorption is so readily distilled out of a more complex situation.

In conclusion, a comparative investigation of laser-light absorption has revealed a very similar behavior at $\lambda = 1.06$ and 0.53 μm. Absorption is dominated by resonance absorption; we have verified its three main characteristics: (i) angle and polarization dependence, (ii) independence of intensity, and (iii) independence of wavelength. Resonance absorption is thus found to be an important mechanism at infrared *and* visible wavelengths.

The authors acknowledge the skillful technical assistance of P. Sachsenmaier and E. Wanka. This work was supported in part by the Bundesministerium für Forschung und Technologie and EURATOM. One of us (A.G.M.M.) acknowledges receipt of a EURATOM grant.

[1]J. S. Pearlman and M. K. Matzen, Phys. Rev. Lett. 39, 140 (1977).

[2]K. R. Manes, V. C. Rupert, J. M. Auerbach, P. Lee, and J. E. Swain, Phys. Rev. Lett. 39, 281 (1977).

[3]J. E. Balmer and T. P. Donaldson, Phys. Rev. Lett. 39, 1084 (1977).

[4]R. P. Godwin, P. Sachsenmaier, and R. Sigel, Phys. Rev. Lett. 39, 1198 (1977).

[5]A. G. M. Maaswinkel, K. Eidmann, and R. Sigel, Max-Planck-Gesellschaft zur Förderung der Wissenschaften eV. Report No. PLF-10/1978 (unpublished).

[6]See, for example, the discussion of experiments by B. H. Ripin, Naval Research Laboratory Memorandum Report No. 3684, 1977 (unpublished).

[7]F. Amiranoff, R. Benattar, R. Fabbro, E. Fabre, C. Garban, C. Popovics, A. Poquerusse, R. Sigel, C. Stenz, J. Virmont, and M. Weinfield, in *Proceedings of the Seventh International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Innsbruck, Austria, 1978* (International Atomic Energy Agency, Vienna, Austria, 1978), paper IAEA-CN-37-D-4.

[8]V. L. Ginzburg, *The Propagation of Electromagnetic Waves in Plasmas* (Pergamon, Oxford, 1970), p. 267.
[9]K. G. Estabrook, E. J. Valeo, and W. L. Kruer, Phys. Fluids 18, 1151 (1975).
[10]J. W. Shearer, Phys. Fluids 14, 183 (1971).
[11]T. W. Johnston and J. M. Dawson, Phys. Fluids 16, 722 (1973).
[12]J. J. Thomson, W. L. Kruer, A. Bruce Langdon, Claire Ellen Max, and W. C. Mead, Phys. Fluids 21, 707 (1978).
[13]Julian Crowell and R. H. Ritchie, J. Opt. Soc. Am. 60, 794 (1970).

Goodness of Ergodic Adiabatic Invariants

Edward Ott
Laboratory of Plasma Studies and Department of Electrical Engineering, Cornell University, Ithaca, New York 14853
(Received 15 March 1979)

For a "slowly" time-dependent Hamiltonian system exhibiting ergodic motion, the volume inside the hypersurface on which the Hamiltonian equals a constant is an adiabatic invariant. It is shown that the error in the constant is diffusive and scales as $(\tau_C/\tau)^{1/2}$, where τ_C is a certain correlation time of the ergodic motion, and τ is the time scale over which the Hamiltonian changes.

The importance of ergodically wandering solutions of Hamilton's equations has been demonstrated in a variety of plasma-physics problems.[1,2] Recently, Lovelace[2] has considered the compression of a field-reversed ion ring. After assuming that the motion of the ring ions should be ergodic in a plane transverse to the toroidal direction, he demonstrated an adiabatic invariant for the ergodically moving ion. Although his derivation was specific to the ion-ring problem, the invariant is actually a very general one. Namely, for a "slowly" time-dependent Hamiltonian system exhibiting ergodic motion in N spatial dimensions $(\vec{q})$, the volume of $2N$-dimensional phase space $(\vec{q},\vec{p})$ within the hypersurface $H(\vec{p},\vec{q},t)=\text{const}$ (where H is the Hamiltonian) is an adiabatic invariant. (Indeed the existence of this adiabatic invariant is already appreciated in statistical mechanics,[3] and the invariant may be associated with the system entropy. In statistical mechanics N is large, whereas $N=2$ is of interest for Refs. 1 and 2.) The generality of the ergodic invariant suggests that it may be useful in a wide variety of other plasma-physics problems where ergodic particle motion is prevalent. Motivated by this, the present work attempts to evaluate the goodness of the ergodic adiabatic invariant. That is, since the adiabatic invariant is only approximately conserved, how good is the approximation? For the case of the familiar $N=1$ adiabatic invariant $(\oint p\,dq)$ of a particle exhibiting rapid almost periodic motion (e.g., the magnetic moment), the average error in assuming that $\oint p\,dq$ is conserved can be exponentially small in τ, where τ is the time scale over which the Hamiltonian changes.[4] In contrast, it is shown here that, for the ergodic invariant, the error is typically proportional to $\tau^{-1/2}$, and it is shown how to calculate the error.

In order to present a brief heuristic demonstration of the ergodic adiabatic invariant, suppose that the existence of three widely separated time scales, $\tau \gg T \gg \tau_w$, where τ_w is the time it takes the system to wander over the surface $H=\text{const}$, where H is the Hamiltonian, and in computing τ_w one uses the orbit obtained from Hamilton's equations with the explicit slow time dependence of H neglected (since τ is large). The exact distribution function of the system is $f=\delta(\vec{p}-\vec{p}(t))\,\delta(\vec{q}-\vec{q}(t))$, where $\vec{p}(t)$ and $\vec{q}(t)$ are solutions of the exact equations of motion. According to the ergodic theorem,

$$\langle f\rangle_T \cong K(t)\delta(H(\vec{p},\vec{q},t)-H_0(t)), \qquad (1)$$

where $\langle\cdots\rangle_T$ denotes an average over the time scale T. Note that K and H_0 evolve on the slow time scale τ (as does H). (For $\partial H/\partial t=0$ the Hamiltonian is a constant of the motion and H_0 is just a constant, but $\partial H/\partial t\neq 0$ is of interest here.) The principal use of the ergodic invariant is that it will determine the time dependence of H_0. If a surface in $(\vec{p},\vec{q})$ phase space is evolved (with each point on the surface following a system orbit), then the volume inside that surface is conserved.[5] Since (1) represents a distribution function, the surface $H=H_0$ evolves in this manner.

