Laser-plasma detection

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When intense biharmonic laser radiation illuminates a charged metal target, the current flowing from the plasma to the target is found to contain an rf component at the difference frequency ~ 0.6 GHz of the laser radiation (this is called laser-plasma detection). It is shown that the expansion velocity of the conducting region of the plasma near the charged metal target is modulated as a function of time, and this makes laser-plasma detection possible for biharmonic Nd lasers with pulse length $\tau \sim 50$ ns. The expansion occurs because the optical breakdown plasma near the surface emits hard ultraviolet radiation which sets up a photoionization wave in the surrounding cold gas.

INTRODUCTION

Laser light irradiating a charged metal surface can induce surface currents and cause a current to flow in an electrical system connected to the target. If the light intensity lies below the threshold for plasma formation, the induced current is due to photo- or thermal emission of electrons from the target surface.[1] Studies have shown[1] that the magnitude of the current pulse depends on the magnitude and sign of the target potential U—the electric signal increases with $|U|$ for $U<0$, but there is no signal for $U \geqslant 0$.

The physical mechanism for current generation is qualitatively different if the laser light is intense enough to form a conducting plasma near the target surface.[2] In this case, static charge on the target flows into the laser-induced plasma generated by surface optical breakdown. The experiments we reported in Ref. 2 show that the sign of the target charge determines the polarity of the current but not its magnitude. We also established that during a nanosecond laser pulse, the magnitude of the current from the target is proportional to the laser light intensity, i.e., the current pulse in the target-plasma system has the same form as the light pulse (Fig. 1). In Ref. 3, we exploited this property to generate ultra-high-frequency (UHF) currents during optical breakdown at the surface of a charged metal target irradiated by laser radiation of modulated intensity. This phenomenon was called laser-plasma detection by analogy with radiowave detection methods in which the envelope of the carrier frequency is selected. In this paper we report experimental results on the physical mechanism for rf current generation.

1. TECHNIQUE FOR GENERATING INTENSE rf-MODULATED LASER BEAMS

The intensity of laser radiation incident on the surface of a target can be modulated at a frequency ν_- by using two light beams of constant intensity and frequencies ν_1, ν_2 to illuminate the target. To achieve a large modulaton factor one must ensure that the two beams have a high degree of spatial and temporal coherence.

The spectral widths $\Delta\nu_1, \Delta\nu_2$ of the light fields determine the coherence times for each beam: $T_1 \sim 1/\Delta\nu_1$, $T_2 \sim 1/\Delta\nu_2$. When the two fields are mixed, periodic modulation of the intensity is observed during a time $T \leqslant \min\{T_1, T_2\}$. To achieve intensity modulation over times $T \sim 10$ nm the spectral width for each beam must be $\Delta\nu \sim 10^{-3}\ \text{cm}^{-1}$.

The current flowing in the plasma-target system depends on the time behavior of the plasma volume, which in turn depends on the energy flux entering the plasma. The energy flux must thus be modulated as a function of time if rf currents are to be obtained. The entire energy flux incident on the target can be modulated if the phase difference of the waves reaching the target is the same at all points on the surface. This condition requires that the wavefronts of the two light beams be identical.

The above requirements can be met by employing an optical system with a frequency-shifted narrow-band pumping source based on stimulated Brillouin scattering (SBS) in active materials.[4,5] The two unequal frequencies are generated by SBS in two materials with different hypersound velocities, and the phase conjugation accompanying the SBS ensures good spatial coherence. Indeed, for SBS backscattering the Stokes shift is $\nu_p - \nu_s = \nu_p(2mv/c)$, where ν_p and ν_s are the pump and the Stokes frequencies, m is the refractive index of the active material, v is the speed of sound, and c is the speed of light in vacuum. Two different SBS materials and a single pumping source of frequency ν_p suffice to generate light fields of frequency $\nu_1 = \nu_p(1 - 2m_1v_1/c)$ and $\nu_2 = \nu_p(1 - 2m_2v_2/c)$, and the intensity of the laser radiation is correspondingly modulated at the frequency $\nu_- = 2c^{-1}|m_1v_1 - m_2v_2|\nu_p$. Frequencies in the range 10^8–$10^{10}\ \text{s}^{-1}$ are possible by combining different active materials and pumping sources in such an optical system.

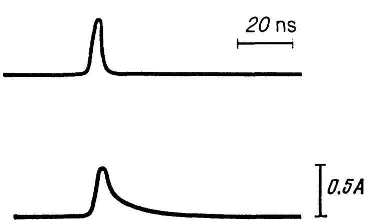

FIG. 1. Radiation pulseform for an Nd laser (top) and trace of the current in the plasma-target system (bottom).

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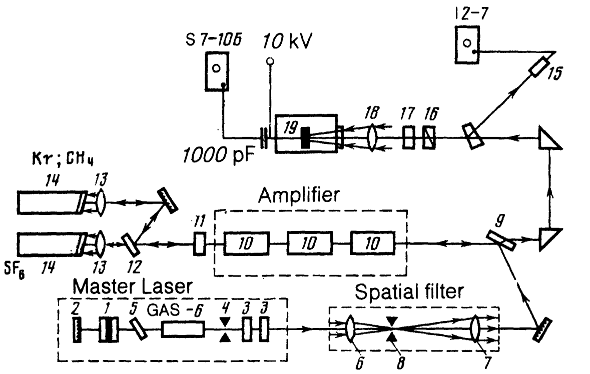

FIG. 2. Schematic of the experimental system.

2. OPTICAL SYSTEM

The optical system used in the experiment (Fig. 2) contained a master laser, a three-stage amplifier with "Brillouin" mirrors, and recording equipment. The active element of the master laser was a 12×270 mm GLS-6 neodymium silicate glass rod, and the laser was Q-switched by a bleachable dye 1. The resonator was bounded by a flat dielectric mirror 2 ($R = 100\%$ at $\lambda = 1.06\,\mu$m) and by two plane-parallel K-8 glass plates 3.

We were able to select the fundamental transverse mode and one longitudinal mode by employing a suitable (complex) resonator configuration and inserting a diaphragm 4 of diameter 3 mm into the resonator. Linear polarization was achieved by mounting a plane-parallel K-8 glass plate 5 in the resonator at the Brewster angle with respect to the optical axis. The master laser generated pulses of length $\tau \sim 70$ ns (half-intensity) and energy $E \sim 20$ mJ, and the spectral width was $\Delta\nu_p \lesssim 10^{-2}\,\mathrm{cm}^{-1}$. Prior to amplification, the light from the master laser passed through a spatial filter consisting of two lenses 6,7 (focal lengths $F_1 = 0.5$ and $F_2 = 1$ m) in air and a diaphragm 8 of diameter $d = 0.5$ mm located at the focus of the first lens. A collimated light beam of diameter 6 mm with a smooth cross sectional intensity distribution was formed at the output of the spatial filter, which also served to decouple the master oscillator from the amplifier.

Roughly 15% of the pumping radiation was input to the amplifying system by reflection from one face of a plane-parallel glass plate 9 mounted at the Brewster angle in a plane parallel to the polarization vector of the pumping radiation. There were three stages, each consisting of a GOR-300 amplifier. The active elements were GLS-22 neodymium phosphate glass rods 10 which measured 12×270 mm; the gain per pass through the entire amplifier system was adjustable from 10^2 to $3 \cdot 10^2$.

The amplified radiation passed through a K-8 glass Fresnel rhombus 11 at the exit from the amplifying system, was split by a mirror 12 ($R = 50\%$) into two beams of equal intensity, and was focused by lenses 13,14 ($F_1 = F_2 = 40$ cm) into one or more 70-cm-long cells 14 filled with compressed SF_6, CH_4, and Kr at pressures 21, 58, and 33 atm, respectively. The backscattered SBS beams were combined at the mirror ($R = 50\%$), passed through the Fresnel rhombus (thereby becoming linearly polarized perpendicular to the polarization of the pumping radiation), and were amplified. Since the reflected radiation was polarized normal to the pumping radiation, there was no reflection from the Brewster glass plate at the entrance to the amplifier, so that the latter was decoupled from the master laser.

Figure 3(a) shows the oscilloscope trace of the Stokes component of the SBS pulse in SF_6; similar traces were recorded for SBS in Kr and CH_4. The reflection coefficient for SBS in SF_6, KR, and CH_4 was $R \sim 5\%$ in our experiments. The SBS in each cell was accompanied by phase conjugation for the pump energies used in the optical system, and this ensured that the two light beams of frequency ν_1, ν_2 were spatially coherent.

Figure 3b,c shows traces of the radiation intensity at the output of the amplifier when the Stokes components were generated in two cells simultaneously. The period of the intensity oscillations was equal to 5 ns for simultaneous SBS in SF_6 and Kr ($v_{SF_6} = 113$ m/s, $v_{Kr} = 224$ m/s) and 1.5 ns for SBS in SF_6 and CH_4 ($v_{CH_4} = 430$ m/s). The laser and plasma emission pulseforms were recorded by coaxial FÉK-15 and FÉK-09 photocells 15. The pumping radiation passed through a polarizer 16 and $\lambda/4$-plate 17 and was focused by a lens 18 ($F = 0.4$ m) onto the positively charged target 19 in an evacuated dielectric chamber.

3. ANALYSIS OF THE CURRENT IN THE PLASMA-TARGET SYSTEM

Figure 4 shows the electrical system used to record the current; $C_d = 1000$ pF is a capacitive divider, $R = 75\,\Omega$ is the input impedance of the oscilloscope, and the capacitance of the target (a 2-mm-thick brass disk of diameter 10 mm) was $C_t \approx 10$ pF, including the leads.

Before each experiment we charged C_d and C_t to a potential difference U_0. When the laser beam struck the target, the surface plasma expanded, C_t increased with time, and the redistribution of charge on the capacitors C_d and $C_t(t)$

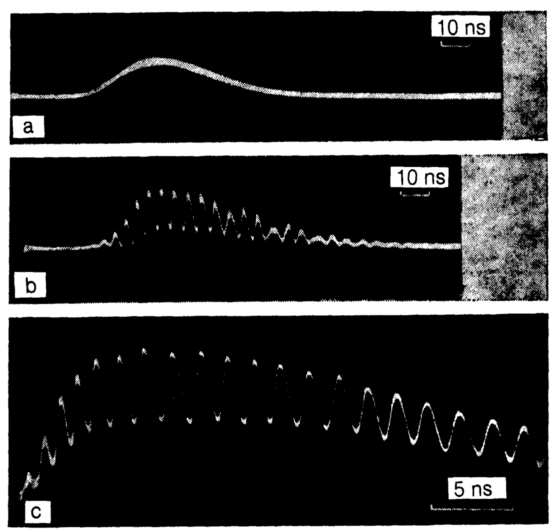

FIG. 3. Radiation pulseforms for SBS in SF_6 (a), $SF_6 + Kr$ (b), and $SF_6 + CH_4$ (c).

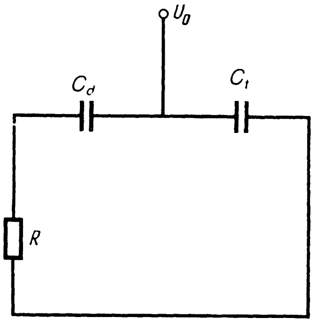

FIG. 4. Circuit used to record the currents.

caused a current to flow in the circuit.

The magnitude of the current can be found by solving the system of equations

$$C_d U_d(t) + C_t(t) U_t(t) = [C_d + C_t(0)] U_0,$$
$$i(t)R + U_d(t) - U_t(t) = 0, \quad dU_d(t)/dt = i(t)/C_d, \tag{1}$$

where $U_d(t)$ and $U_t(t)$ are the voltages across C_d and $C_t(t)$ and $i(t)$ is the current. If $C_t(t)$ increases linearly with time, $C_t(t) = C_t^0 + Vt$, system (1) has the solution[1]

$$i(t) = \frac{U_0 V}{VR+1}\left[\left(\frac{C_t^0}{C_t^0+Vt}\right)^{(VR+1)/VR} - 1\right]. \tag{2}$$

If $VR \gg 1$ (so that C_t increases rapidly),

$$i(t) \approx \frac{U_0}{R}\left[\left(\frac{C_t^0}{C_t^0+Vt}\right) - 1\right] \xrightarrow[t\to\infty]{} -\frac{U_0}{R}. \tag{3}$$

In the opposite limit $VR \ll 1$, the current is given by

$$i(t) \approx U_0 V\left[\left(\frac{C_t^0}{C_t^0+Vt}\right)^{1/VR} - 1\right] \xrightarrow[t\to\infty]{} -U_0 V. \tag{4}$$

For small times $Vt \ll C_t^0$, expression (4) can be rewritten as

$$i(t) \approx U_0 V[\exp(-t/RC_t^0) - 1], \tag{5}$$

which gives $\tau = RC_t^0$ for the relaxation time constant when C_t increases slowly.

The capacitance $C_t(t)$ of the target is equal to the sum

$$C_t(t) = C_{md}(t) + C_{cp},$$

where the capacitance C_c of the leads is independent of t; C_{md} is the intrinsic capacitance of the metal disk, $C_{md} = 8\varepsilon_0 r(1 - D(t)/\pi r)^{-1}$ (Ref. 6), where r and $D(t)$ are the radius and thickness of the disk.

The thickness of the disk (assumed to be flat) increases during the formation of the highly conductive laser-induced surface plasma. The rate V at which the target capacitance increases is thus related to the rate v of plasma expansion near the surface of the disk:

$$V = \frac{dC_t(t)}{dt} = \frac{dC_{md}(t)}{dt}$$
$$\approx \frac{8\varepsilon_0}{\pi}\frac{dD}{dt} = \frac{8\varepsilon_0}{\pi} v. \tag{6}$$

This expression is valid when $D(t) \ll r$.

For our system, $R = 75\ \Omega$ and the inequality $VR \ll 1$ holds for all $0 \leqslant v \leqslant c$; $i(t)$ is therefore given by (4). It follows that if v is modulated in time with a period T less than the relaxation time τ ($\tau \sim 1$ ns for our system), $i(t)$ should contain an rf component with the same period T.

4. PLASMA EXPANSION MECHANISMS

Rf currents of period T are generated in the plasma-target system when the plasma expansion rate v is time-modulated and the charge in the plasma is able to redistribute itself during a time $t < T$, so that the potential in the conducting plasma region becomes equalized. The charge redistribution time is determined by the Maxwell relaxation time,[2] which is equal to[7]

$$T_M = 1/4\pi\sigma, \tag{7}$$

where the plasma conductivity σ is given by

$$\sigma = ne^2/\nu_{ea} m_e \tag{8}$$

(n is the density of free electrons, ν_{ea} is the free electron-atom collision frequency, and m_e is the electron mass).

We can use this relation to find the conducting region of the plasma in which modulation of v with period T will produce a corresponding modulation of the current in the plasma-target system; it is defined by the condition $n \geqslant n_{cr}$, where

$$n_{cr} = \nu_{ea} m_e / 4\pi e^2 T. \tag{9}$$

The conducting region that forms near the target surface during low-threshold optical breakdown can be divided into two parts—a region of dense ($n \sim 10^{18}$–10^{19} cm^{-3}) highly conducting plasma which absorbs laser radiation,[8] and a relatively poorly conducting region which is produced by photoionization of the gas by the hard ultraviolet radiation emitted by the dense plasma.[9,10] We consider these two regions separately.

1. Under our experimental conditions (laser wavelength $\lambda = 1.06\ \mu$m, intensity $I \sim 10^8$–10^9 W/cm^2), a light detonation wave may be generated in the plasma due to low-threshold optical breakdown at the surface. In this case the laser radiation is absorbed at the leading edge of a shock wave that travels through the gas surrounding the target. The electron density behind the detonation wave is $n_e \sim 10^{18}$–10^{19} cm^{-3} when the gas pressure ahead of the wavefront is $P \sim 1$ atm, and the temperature T is ~ 1–2 eV (Ref. 8). The velocity of the light detonation wave depends on the laser intensity and is given by[11]

$$D = [2(\gamma^2 - 1) I_l/\rho]^{1/3}, \tag{10}$$

where γ is the adiabatic exponent of the gas, I_l is the laser intensity, and ρ is the gas density ahead of the wavefront.

We see from (10) that the velocity of the detonation wave (and hence also the expansion velocity v of the conducting region) can be modulated by modulating the laser radiation as a function of time. For $I_l \sim 10^8$ W/cm^2, we have $D \sim 10^6$ cm/s.

2. In the poorly conducting region, the UV radiation from the dense plasma region ionizes the cold gas (air) surrounding the plasma. Ions and free electrons accumulated

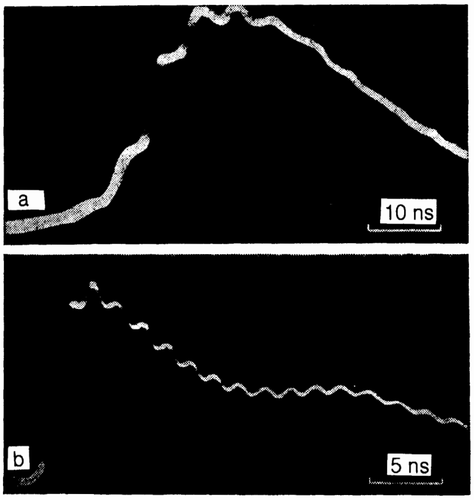

FIG. 5. Traces of the current pulses in the target-plasma system; the laser radiation was modulated with period $T \approx 5$ ns ($\nu \sim 200$ MHz) (a) and $T \approx 1.5$ ns ($\nu \sim 600$ MHz) (b).

much faster than they recombine, so that the conducting region expands. If we let τ_e be the lifetime of a free electron in a neutral gas, i.e., the time required for attachment to a neutral atom (τ_e is $\sim 10^{-7}$ s (Ref. 12) for air at $P = 1$ atm) and consider the UV ionization for times $t < \tau_e$, then the condition $n = n_{cr}$ defines the boundary of the expanding region at a given time and given modulation frequency.

Let us examine how the ionization wave moves in the one-dimensional case. We assume that the plasma is ionized by UV radiation of frequency ν and write α for the UV absorption coefficient in the cold gas; the plasma emission intensity $I_{UV}(t)$ at the frequency ν is assumed to be specified. We choose the z axis so that the surface of the bright region (dense plasma) lies in the plane $z = 0$, while the radiation flux is directed along the positive z axis. In addition, we assume that only one photon is required for ionization. The free electron density at time t and point z is then given by

$$n(z,t) = \frac{1}{h\nu} \int_{z/c}^{t} \alpha I_{UV}\left(\xi - \frac{z}{c}\right) e^{-\alpha z}\, d\xi. \tag{11}$$

If $t \gg z/c$, or if the modulation period T is much greater than the time required for the light to pass through the conducting region, (11) simplifies to

$$n(z,t) \approx \frac{1}{h\nu} \int_{0}^{t} \alpha I_{UV}(\xi)\, e^{-\alpha z}\, d\xi. \tag{12}$$

We can use this to find the velocity v of a point corresponding to a specified density n (e.g., $n = n_{cr}$) along the z axis; we find that

$$v = \left.\frac{dz}{dt}\right|_{n=\text{const}} = I_{UV}(t) \Big/ \alpha \int_{0}^{t} I_{UV}(\xi)\, d\xi. \tag{13}$$

We assume that the absorption of the modulated laser light in the plasma produces corresponding oscillations in the plasma brightness with the same modulation period T. The plasma emission intensity can then be expressed as

$$I_{UV}(t) = I_0[1 + \delta(t)],$$

where the function $\delta(t)$ is periodic with period T. If $|\delta(t)| \ll 1$, we get the following expression for the expansion rate of the conducting region:

$$v \approx [1 + \delta(t)]/\alpha t. \tag{14}$$

This shows that v is modulated, and the oscillations induce an *ac* current with period T in the circuit (Fig. 4).

For $t \sim 10$ ns and absorption $\alpha \sim 10$ cm^{-1}, the rate of expansion due to the photoionization wave can reach $v \sim 10^7$ cm/s, which is an order of magnitude greater than for a light detonation wave. However, (14) implies that v decreases quite rapidly with α and t.

We note that the expansion of the conducting region $n \geqslant n_{cr}$ is not coincident with but rather lags the emission from the plasma, because a finite time is required for the electron density in the gas adjacent to the plasma to reach the critical value. The delay t_d becomes shorter as I_{UV} increases and is given by

$$t_d \approx n_{cr} h\nu / I_{UV} \alpha. \tag{15}$$

5. EXPERIMENTAL RESULTS AND DISCUSSION

Figure 5 shows some traces for the current in the plasma-target system for a target irradiated by intensity-modulated laser pulses at frequency $1/T = 200$ MHz (a) and 600 MHz (b). The target was located at a distance $L = 30$ cm from the lens, the diameter of the focal spot was 2 mm, and the laser intensity was $I_l = 2 \cdot 10^8$ W/cm^2. We see from the figures that the current pulses were amplitude-modulated at the same frequency $1/T$ as the laser radiation, i.e., the current $i(t)$ flowing from the target into the plasma contained an rf component. The modulation persisted when the laser intensity I_l was increased from 10^8 to 10^9 W/cm^2.

Formulas (4),(6) give the estimate $v \sim 10^7$ cm/s for the plasma expansion; this is nearly an order of magnitude greater than the velocity of light detonation waves for these laser intensities. It follows that in our situation, photoionization of the gas by ultraviolet radiation from the plasma (photoionization wave) is the dominant factor in the expansion of the conducting region. We verified this conclusion by analyzing the emission from the laser plasma. Figure 6 shows some plasma emission traces for an intensity-modulated laser beam incident on the target. We see that $I_{UV}(t)$ contains an ac component of frequency $1/T$ equal to the laser modulation frequency. According to (14), the ac component in $I_{UV}(t)$ induces an rf component in the current $i(t)$, as was observed experimentally. Analysis of the plasma emission also reveals that for $0.01 \leqslant P \leqslant 1$ atm, the pressure P of the gas (air) surrounding the target had no influence on the UV plasma emission ($\lambda_{UV} \approx 400$ nm) during the first 50 ns after optical breakdown, as determined from the onset of intense plasma emission. This is because the laser absorption region was unable to expand into the air from the vaporized target during this time, i.e., no light detonation wave formed in the air. Indeed, for intensities $I_l \approx 10^8$ W/cm^2, air pressures $P \approx 1$ atm, and plasma temperatures $T_{pl} \approx 2$ eV, the width of the absorption wavefront is equal to $\sim 10^{-2}$ cm, the photon

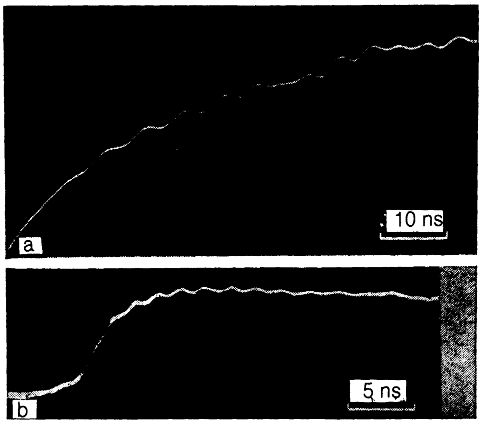

FIG. 6. UV emission from the laser plasma in the 350–450 nm range; the modulation period of the laser radiation was $T \approx 5$ ns ($\nu \sim 200$ MHz) (a) and $T \approx 1.5$ ns ($\nu \sim 600$ MHz) (b).

mean free path for an Nd laser at wavelength $\lambda = 1.06\,\mu$m (Ref. 13); since a light detonation wave moves at velocity $D \sim 10^6$ cm/s, during 10 ns the wave can only travel a distance $\sim 10^{-2}$ cm comparable to the width of the wavefront. The plasma (vaporized target material) is therefore not screened, and the plasma emission is initially independent of the ambient gas pressure.

Since $I_{\rm UV}(t)$ is independent of P under these conditions, (13) shows that the expansion rate of the conducting region, and hence also the current, depends on the gas pressure as $v \propto P^{-1}$ (because $\alpha \propto P$). Figure 7 shows the current amplitude in the plasma-target system as a function of the ambient pressure on a semilogarithmic scale; we see that $i \propto P^{-1}$ to within the experimental error. This dependence also confirms that for the laser parameters considered, the expansion of the conducting region behind the photoionization wave is responsible for the current. We also conclude from the dependence $i(P)$ in Fig. 7 that under the experimental conditions, the expansion was nearly one-dimensional (the transverse diameter of the plasma was equal to the diameter ~ 2 mm of the laser spot on the target). This implies that the characteristic length of the photoionization region was $\leqslant 1$ mm and gives $\alpha \sim 10$–100 cm^{-1} for the UV absorption coefficient ($\lambda_{\rm UV} = 100$ nm) in air, in agreement with tabulated data.[14,15]

The rf modulation of the current persisted longer as the modulation period T increased (Fig. 5) and was equal to 30–40 ns and 15–20 ns for $1/T = 200$ and 600 MHz, respectively. The slower rate of expansion (14) and the drop in the modulation factor for $I_{\rm UV}$ (Fig. 6) are apparently responsible for this behavior.

We attribute the modulation of the plasma emission to a discontinuous change in the temperature of the plasma layer in which the laser radiation is absorbed. We found that during laser-induced breakdown at $I_l \sim 10^{10}$ W/cm^2, $I_{\rm UV}$ was modulated at the same frequency as the laser light when the latter was modulated at $1/T = 200$ MHz. However, the oscillations in $I_{\rm UV}$ disappeared when $1/T$ was increased to 600 MHz; in this case, the diameter of the absorbing region was $l \sim 10^{-8}$ cm (Ref. 8), and the laser energy contained in a single spike modulated at 600 MHz was too low to cause significant heating and change the plasma brightness. The absorption coefficient α at the laser wavelength increases during absorption, so that the mass of the target material heated by the laser radiation decreases and a single laser spike can appreciably heat the plasma. The increase in α is a consequence of the fact that the ionization potential of the metal atoms is less than for the atomic constituents of air, and it can be calculated by the Kramers-Unsold formula[16] $\alpha \sim \exp[-(I - h\nu)/kT]$, where I is the ionization potential, $h\nu$ is the laser photon energy, and T is the plasma temperature. Thus for laser photons of energy ~ 1 eV, the difference between the ionization potentials $I = 7.7$ eV and $I = 12$ eV for Cu and O_2 is reflected in a 50-fold increase in α for a specified plasma density and temperature ($T = 1$ eV).

CONCLUSIONS

In this paper we have thus established the following.

1. Laser-plasma detection (i.e., measurement of the difference frequency) for two Nd laser pulses of length 50 ns is possible because the expansion velocity of the conducting plasma near the charged metal target is modulated in time. Hard UV radiation from the hot dense plasma produced by optical breakdown photoionizes the cold gas and is responsible for the expansion.

2. The UV plasma emission, and hence also the current, is modulated most effectively if the plasma strongly absorbs the laser radiation. We thus conclude that infrared (IR) lasers (e.g., CO_2 lasers with $\lambda = 10\,\mu$m) should be useful in laser-plasma detection because the inverse bremsstrahlung coefficient in the plasma increases with λ. Ultraviolet radiation from a second laser can be used to generate a plasma which is then exposed to intensity-modulated radiation from an IR laser. By employing IR lasers in combination with UV lasers, which have a low threshold intensity for optical breakdown,[17] one can decrease the energy required for laser-plasma detection.

We note that laser-plasma detection should also be observed when UV laser radiation ionizes the gas surrounding a charged target and a conducting plasma region is formed.

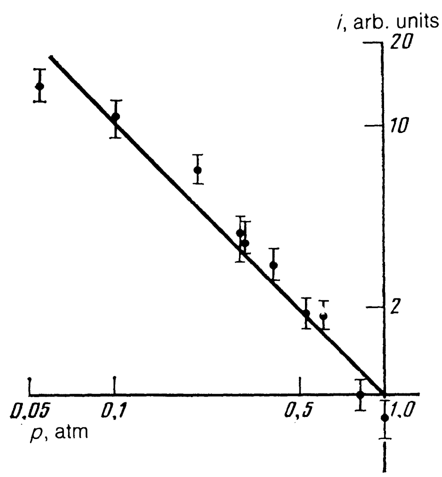

FIG. 7. Current in the plasma-target system as a function of the ambient gas pressure.

[1)]For times $0<t<(C_d - C_t^0)/V$, this is equivalent to the condition $C_t(t) \ll C_d$.

[2)]We assume that the modulation frequency is much less than the frequency of electron collisions with atoms and ions.

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