Ultrashort-pulse laser machining of dielectric materials

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There is a strong deviation from the usual $\tau^{1/2}$ scaling of laser damage fluence for pulses below 10 ps in dielectric materials. This behavior is a result of the transition from a thermally dominated damage mechanism to one dominated by plasma formation on a time scale too short for significant energy transfer to the lattice. This new mechanism of damage (material removal) is accompanied by a qualitative change in the morphology of the interaction site and essentially no collateral damage. High precision machining of all dielectrics (oxides, fluorides, explosives, teeth, glasses, ceramics, SiC, etc.) with no thermal shock or distortion of the remaining material by this mechanism is described. © 1999 American Institute of Physics.

INTRODUCTION

Dielectric materials (e.g., ceramics, SiC, diamond, sapphire, bone, etc.) are very difficult to machine either with abrasive tools or with conventional laser processing. Laser processing by molecular dissociation in organic (and some inorganic) materials can be achieved with ultraviolet lasers (e.g., excimer lasers: KrF, XeCl) but this photodissociation mechanism is not applicable to all dielectrics. By using ultrashort ($\tau<10$ ps) laser pulses any material can be machined to very high precision. The ability to machine any material including high band gap dielectrics such as SiC, diamond, etc. is a result of the fact that these ultrashort pulses interact by a mechanism which is very different from that of conventional longer pulse lasers. This interaction is independent of the usual linear absorption properties of the material and is applicable to materials which would otherwise be transparent to the laser wavelength. Machining to micron scale precision with minimal collateral damage to the remaining material is achieved by removing material faster than heat can be conducted to the bulk.

By dielectrics, we are generally referring to materials with no free electrons and low thermal/electrical conductivity. Common materials such as fused silica, sapphire, SiC, diamond, SiN, AlTiC, ZrO2, glass, plastic, bone, cornea, heart tissue, etc., fall under this category. Semiconductors such as silicon and gallium arsenide would generally not be considered dielectrics by this definition. However, these materials behave similar to dielectrics when machining with ultrashort laser pulses.

Attempts to machine dielectrics with lasers followed shortly after laser induced damage was observed in transparent solids. For pulses longer than a few tens of picoseconds, the generally accepted picture of damage to dielectrics involves the heating of seed electrons by the incident radiation and transfer of this energy to the lattice. Damage occurs via conventional thermal deposition resulting in melting and boiling of the dielectric material. Because the controlling rate is that of thermal conduction through the lattice, this model predicts a $\tau^{1/2}$ dependence of the threshold fluence (energy/area) upon pulse duration $\tau$. This was found to be in reasonably good agreement with numerous experiments which have observed a $\tau^{1/4}$ scaling with $0.4<\alpha<0.5$ in a variety of dielectric materials from 100 ps to microseconds. With these conventional lasers, material is removed by thermal ablation wherein the material is locally heated to near the boiling point. Since the boiling point of these materials is very high (typically >1000 °C), this ablation mechanism is accompanied by a strong thermal shock to the remaining bulk material. This thermal shock often results in cracking of the remaining material and uncontrolled material removal. These effects can be observed in Fig. 1(a) where the initial stages of hole drilling in a tooth using a conventional neodymium laser are shown. Linear absorption due to defects produces inhomogeneous energy absorption across the laser beam. Thermal stresses increase to the point where ablation begins first from the point with the least material strength.

It had been postulated that the $\tau^{1/2}$ dependence of the laser damage threshold for dielectrics would break down for pulse durations less than a few picoseconds as the probability of multiphoton ionization of the individual atoms within the dielectric became significant. A strong deviation from the $\tau^{1/2}$ dependence of the damage threshold was observed by Du et al. in SiO2, and by Stuart et al. in a variety of dielectric materials. This deviation begins at $\sim$10 ps (SiO2) to $\sim$20 ps (CaF2) dependent upon the dielectric (Fig. 2).

In all cases, the deviation was accompanied by a dramatic change in morphology of the damaged surface [Fig. 1(b)]. With femtosecond pulses, all regions throughout the laser beam profile with sufficient intensity for multiphoton ionization will be removed resulting in extremely fine control of the position of material removal. In teeth, the morphology of the surface drilled with femtosecond pulses is characteristic of the internal enamel and dentin. There is no evidence of heat transfer into the surrounding material and no thermal shock induced cracking [Fig. 1(c)].
ond pulses faster than it is transferred from the electrons to the lattice. For these very short, intense pulses, energy is produces free electrons which are then rapidly accelerated by the laser pulse. Field-induced multiphoton ionization processes thereby producing more free electrons. This process continues until a critical density plasma is reached wherein minimal further energy deposition from the laser occurs. The actual damage occurs after the pulse has passed, when the dense plasma expands away from the surface.

Plasma formation is quantitatively described by the time dependence of the electron energy distribution function. Since the impact ionization rate is energy dependent, the absorption rate of laser energy requires integrating over the electron energy distribution. When the electrons are strongly driven by intense laser pulses, the energy distribution can differ substantially from a Maxwellian. For material having a band gap energy which is much larger than the single photon energy \( U_p = h\omega \), the heating and collisional ionization of conduction electrons can be described by a Fokker–Planck equation\(^5,12\) for the electron distribution function \( N(\epsilon,t) \),

\[
\frac{\partial N(\epsilon,t)}{\partial t} + \frac{\partial}{\partial \epsilon} \left( R_J(\epsilon,t) N(\epsilon,t) - \gamma(\epsilon) E_p N(\epsilon,t) \right) - D(\epsilon,t) \frac{\partial N(\epsilon,t)}{\partial \epsilon} = S(\epsilon,t),
\]

where \( \epsilon \) is the electron kinetic energy. The number of electrons with a kinetic energy in the range: \( \epsilon \) to \( \epsilon + d\epsilon \) at the time \( t \) is given by \( N(\epsilon,t) d\epsilon \). The bracket in Eq. (1) represents the change in the electron distribution due to Joule heating, inelastic scattering of phonons, and electron energy diffusion. The Joule heating rate of the electrons is given by

\[
R_J(\epsilon,t) = \frac{\sigma(\epsilon)}{3} E_{rms}(t)^2, \tag{2}
\]

where \( E_{rms}(t) \) is the electric field of the laser and \( \sigma(\epsilon) = e^2 \tau_m(\epsilon)/m^* [1 + \omega^2 \tau_m(\epsilon)^2] \) is the ac conductivity of a conduction-band electron of effective mass \( m^* \). Here, \( \tau_m(\epsilon) = 1/\tau_e(\epsilon) \) is the energy-dependent, electron–phonon transport (momentum) scattering rate. It is instructive to rewrite this expression as

\[
R_J(\epsilon,t) = \left( \frac{1}{3} \right) \left( \frac{m_e}{m^*} \right) \nu_{\text{col}}(\epsilon) U_p(t), \tag{3}
\]

where \( U_p = e^2 E(t)^2/4m\omega^2 = 9.33 \times 10^{-14} l/W/cm^2 \) is the free electron quiver energy expressed in eV and \( \nu_{\text{col}}(\epsilon) = \nu_{\text{mp}}(\epsilon) \omega^2 \tau_m(\epsilon)^2/[1 + \omega^2 \tau_m(\epsilon)^2] \approx \nu_{\text{mp}}(\epsilon) \) is the effective electron collision frequency with the lattice. Under the conditions of interest here, the effective mass \( m^* \) is nearly identical to the electron rest mass \( m_e \). For SiO\(_2\), the momentum scattering rate \( \nu_{\text{mp}}(\epsilon) \) varies between 2 and 10 fs\(^{-1}\) from 1 to 10 eV.\(^{13}\) Hence, for a laser wavelength of 1064 nm and an intensity of 10\(^{13}\) W/cm\(^2\), the single electron heating rate is on the order of a few eV/fs.

The rate of energy transfer from the electrons to the lattice is given by the second term within the brackets \( \gamma(\epsilon) E_p(t) N(\epsilon,t) \), where \( E_p \) is the energy of a typical phonon and \( \gamma(\epsilon) \) is the rate of electron–phonon energy transfer to the lattice. For pulses on the order of 1 TW/cm\(^2\), the electron heating rate \( R_J(\epsilon,t) \) is much greater than the rate of energy transfer to the lattice \( \gamma(\epsilon) E_p(t) \approx 0.01 eV/ps \). The term \( D(\epsilon,t) \partial N/\partial \epsilon \) represents electron energy diffusion with a diffusion coefficient given by \( D(\epsilon,t) = 2\epsilon R_J(\epsilon,t) \).
The final term represents sources and sinks of electrons. Specifically, the total number of electrons created with kinetic energy in the range \( e \) to \( e + de \) at the time \( t \) is given by \( S(e,t)de \). Free electrons can be created by both multiphoton ionization and collisional (impact) ionization. Electrons which have been accelerated to energies above the ionization potential \( U_{IP} \) can collisionally ionize a neighboring atom. An incident electron of kinetic energy \( e_0 = 2e + U_{IP} \) produces two free electrons of energy: \( e \) and an ion (hole) of potential energy \( U_{IP} \) in the final state. Impact ionization occurs at a rate described by the Keldysh impact formula,

\[
\nu_I(e_0) = \chi(e_0/U_{IP} - 1)^2. \tag{14}
\]

For fused silica, the proportionality constant \( \chi = 1.5 \times 10^{15} \text{s}^{-1} \) (Ref. 13) for incident electrons of energy \( < 100 \text{ eV} \).

The source term for impact ionization can be written as

\[
S(e,t)\text{impact} de = 2N(e_0,t)\nu_I(e_0)d\epsilon - N(e,t)\nu_I(e)de = [4N(2e + U_{IP},t)\nu_I(2e + U_{IP}) - N(e,t)\nu_I(e)]de. \tag{4}
\]

The term, \( 4N(2e + U_{IP},t)\nu_I(2e + U_{IP}) \) represents those electrons born into \( e \) to \( e + de \) by impact ionization, while the term \( N(e,t)\nu_I(e) \) represents those lost from this energy bin. The factor of 4 in the source term can be understood simply from the fact that the two electrons born into \( e \) and \( e + de \) from the initial energy \( e_0 = 2e + U_{IP} \) are born into an energy bin \( de \), which is only half as large as the original (e.g., \( d\epsilon_0 = 2de \)).

The total source term is given by

\[
S(e,t)de = [S(e,t)\text{pi} + S(e,t)\text{impact}]de = [S(e,t)\text{pi} + 4N(2e + U_{IP},t)\nu_I(2e + U_{IP}) - N(e,t)\nu_I(e)]de. \tag{5}
\]

The photoionization source term \( S(e,t)\text{pi}de \) has been studied extensively in the case of isolated atoms. In the limit that the band gap of the material is not too much greater than the photon energy \( U_{IP}/h\nu < 4 \), and there are no intermediate resonances, simple perturbation theory can be applied:

\[
\nu_I(e_0) = a^{(K)} e^{K} \text{exp}(-e/K)\text{exp}(e/h\nu). \tag{15}
\]

The energy of the resulting electron is \( e = K\hbar\nu - U_{IP} \). An example of this case would be fused silica irradiated by 532 nm light. Four 2.33 eV photons are required to ionize an SiO2 molecule (\( U_{IP} = 9 \text{ eV} \)) with a cross section similar to that measured for NaCl, \( a^{(K)} = 2 \times 10^{-11} \text{cm}^2 \text{s}^{-1} \).\(^2\) If there are intermediate multiphoton resonances involved, more sophisticated expressions for the ionization rate are required.\(^15\)

In the limit that the IP is less than approximately twice the free electron quiver energy, \( U_{IP} < 2U_F \), perturbation theory breaks down. In this case, there are a large number of photons absorbed (typically \( > 6 \)) and higher order processes such as above-threshold ionization contribute. Photoionization is then described by the tunneling theory of Keldysh\(^16\) or derivatives thereof.\(^17,18\)

Direct numerical solution of Eq. (1) can be found in Ref. 19. However, a very useful approximate solution can be obtained for the total electron density \( N(t) \) in the limit that the electron heating rate is much greater than the rate of energy transfer to the lattice. Note that

\[
\frac{dN(t)}{dt} = \frac{\partial}{\partial t} \int_0^\infty N(e,t)de. \tag{6}
\]

Inserting Eqs. (2)–(5) into Eq. (1) and integrating over energy, we can write the time dependence of the electron density as

\[
\frac{dN(t)}{dt} = \beta(I)N(t) + P(I), \tag{7}
\]

where \( \beta = \alpha I(t) \) contains all the terms associated with electron heating and impact ionization and \( P(I) \) is the photoionization term,

\[
P(I) = \int_0^\infty S(e,t)\text{pi}de = N_0 W(t), \tag{8}
\]

where the term \( N_0 \) is the initial atom density and \( W(t) \) is the single atom photoionization rate. Equation (7) is readily solved to yield

\[
N(t) = N_0 \text{exp} \left[ -\alpha \int_0^t I(t')dt' \right] \int_{-\infty}^t W(t') \text{exp} \left[ -\alpha \int_{t'}^t I(t'')dt'' \right] dt'. \tag{9}
\]

The avalanche coefficient, \( \alpha = 0.01 \text{ cm}^2 \text{s}^{-1} \text{GW}^{-1} \) for fused silica is found by numerically solving Eq. (1)\(^9\) or by fitting the experimental data. Note that \( 1/\alpha = 0.1 \text{ J/cm}^2 \) has the units of fluence.

Many mode-locked short-pulse lasers exhibit a temporal distribution given by

\[
I(t) = I_0 \text{sech}^2(2t/\tau_p). \tag{10}
\]

Inserting this expression into Eq. (9) yields

\[
N(t) = N_0 \text{exp} \left[ -\alpha I_0 \tau_p/2 \text{tanh}(2t/\tau_p) \right] \int_{-\infty}^t W(t') \text{exp} \left[ -\alpha I_0 \tau_p/2 \text{tanh}(2t'/\tau_p) \right] dt'. \tag{11}
\]

This is a very useful analytic expression which predicts the time-dependent electron density from both multiphoton and collisional ionization given an expression for the single-atom multiphoton ionization rate \( W(t) \).

Figure 3 shows the evolution of electron density for a 10 TW/cm\(^2\), 100 fs pulse incident on fused silica. The temporal profile of the laser pulse and the electron density produced by multiphoton ionization alone are included for reference. Because multiphoton ionization is strongly intensity dependent the electron production takes place principally at the peak of the pulse. For this 100 fs duration, multiphoton ionization produces a substantial amount of free electrons. When the electron density produced by multiphoton ionization approaches \( \approx 10^{17} \text{ cm}^{-3} \), the collisional ionization rate begins to exceed the multiphoton ionization rate. Once a high free electron density is produced by multiphoton ioniza-
tion the material no longer has the properties of a dielectric. It is now a conductor and will absorb the laser energy via inverse Bremsstrahlung (Joule) heating similar to a metal or plasma. It is for this reason that both dielectrics and metals have similar behavior and morphology when machined with ultrashort pulses. In essence, the dielectric is converted to a metallic state within the first few tens of femtoseconds. As the laser intensity decreases past the peak of the pulse, the driving force for both multiphoton and collisional ionization ceases and the free electron density remains constant.

For laser pulses of different duration, the relative fraction of multiphoton ionization to avalanche ionization will change. Multiphoton ionization will contribute a relatively greater fraction of the electron density with shorter pulses while avalanche ionization will dominate with longer pulses. In all cases, the electron density saturates at the critical density \(N_c\). The critical density is the density at which the plasma frequency, \(\omega_{pe}(\text{rad/s}) = (4 \pi N_c e^2/m)^{1/2} = (2.21 \times 10^3 N_c (\text{cm}^{-3}))^{1/2}\), is equal to the laser frequency, \(\omega = 2 \pi c/\lambda\). The critical density may be written as \(N_c = (\pi m_e c^2 e^2 \lambda^2)\). At a laser wavelength of 1064 nm, \(N_c = 0.98 \times 10^{21} \text{ cm}^{-3}\). Note that the ratio \(N_c/N_e = \omega_{pe}^2/\omega^2\).

The critical density plays an extremely important role in the interaction of electromagnetic waves with plasmas. This can best be illustrated by considering the dielectric function of the plasma

\[
\varepsilon(\omega) = 1 - N_c \frac{\sigma(\omega)}{\omega} (1 + \omega \tau_n) = 1 - \frac{\omega_{pe}^2}{\omega(\omega + i\nu_n)}.
\]

The dielectric function determines both the refractive index and absorption of electromagnetic waves within the plasma. With a sharp gradient, absorption and reflection can be estimated from the Fresnel formulas. For normal incidence, the absorption, reflection, and transmission coefficients, \(A\), \(R\), and \(T\), respectively, are given by

\[
R = \frac{4 \text{Re} \frac{\varepsilon}{\varepsilon + 1} \varepsilon}{1 + \varepsilon} \quad T = \frac{2 \text{Re} \frac{\varepsilon}{\varepsilon + 1} \varepsilon}{1 + \varepsilon} \quad A = 1 - [1 - \frac{\varepsilon}{\varepsilon + 1}] \approx 4 \text{Re} \left(\frac{\varepsilon}{\varepsilon + 1}\right) \approx 2 \frac{\nu}{\omega} \left(\frac{N_c}{N_e}\right)^{1/2}.
\]

The approximate expression for the absorption is only applicable to low density plasmas where \(\nu < \omega\). At the critical density, the dielectric function becomes 

\[
\varepsilon_c(\omega) = 1 - 1/[(1 + i\nu_m/\omega)].
\]

For the initial irradiation when the temperature of the plasma is low, there are no reliable theoretical estimates of the collision frequency and we use it as an adjustable parameter to match experimental data. At higher temperatures, the Spitzer formula for the electron collision rate can be used

\[
\nu = n_e Z \pi e^4 \ln \Lambda / [(4 \pi e^2)/(m_e kT)^{3/2}] \\
\approx 2.7 \times 10^{-6} n_e Z \ln \Lambda / \Lambda [T_e^{1/2} (s^{-1})],
\]

where \(T\) is the temperature in eV and \(\Lambda\) is the Coulomb logarithm. In near solid density plasmas, this formula is applicable only for temperatures over \(\approx 50\) eV. At lower temperatures, the number of particles per Debye sphere is small and the plasma is nonideal with the collision rate comparable to the plasma frequency.

Formation of a critical density plasma has an important effect in materials processing with short pulse lasers. Once the critical surface is formed, laser energy is either absorbed at that surface or reflected from it. For most plasmas, there is strong absorption at the critical surface since the scale length \(L\) is equal to or greater than the wavelength of the incident light \(L > \lambda\). The maximum scale length is approximately

\[
L_{\text{max}} \approx v_s \tau_p,
\]

where the ion sound velocity \(v_s = (ZkT_e/m_{\text{ion}})^{1/2} \approx 1 \times 10^6 \text{ cm/s}\). Since the pulses of interest here are \(\approx 10\) ps, the maximum scale length is \(L_{\text{max}} \approx 0.1 \mu m\) which is much less than the incident wavelength. In the limit that \(L < \lambda\), \(A \sim \nu/\omega \approx 1\), i.e., the absorption is weak in all but the most collisional plasmas and most of the
electromagnetic energy is reflected from the critical surface. In Fig. 4, we show the reflectivity of a normally transparent fused silica surface as a function of incident laser intensity. Near the damage threshold ($\approx 10^{13} \text{ W/cm}^2$), critical density is not produced until late in the pulse (Fig. 3). Only the last part of the laser pulse will experience any strong reflection. However, by operating high above threshold, critical density is achieved early in the pulse and reflectivities exceeding 90% can be obtained. A measurement of the phase front of the reflected waves shows that the maximum scale length in the plasma over the irradiance range investigated was $<0.1 \lambda$.

The damage threshold of fused silica calculated according to the plasma model just described is compared to measured damage thresholds at both 526 and 1053 nm from 10 fs to over 10 ps in Fig. 5. The theoretical damage fluence defining damage as the point at which the free electron density reaches the plasma critical density ($\approx 10^{21} \text{ cm}^{-3}$) is calculated with no adjustable parameters (Fig. 5, solid lines). Due to the rapid avalanche following production of the seed electrons by multiphoton ionization, the predicted damage threshold is only weakly dependent upon the actual free electron density at which damage occurs. A lower limit would correspond to the condition that the energy density of the conduction electrons equals the binding energy of the lattice ($\approx 10^{19} \text{ cm}^{-3}$). Use of this lower limit serves only to reduce the predicted damage threshold by $\approx 20\%$. In addition to the good agreement with the measurements presented here, experiments by Auderbert,22 von der Linde,23 Krausz,24 and Kautek25 are also in good agreement with our calculations but not those of Du.9 For pulses less than 100 fs, the predicted damage threshold asymptotically approaches the multiphoton limit. In the long-pulse regime ($\tau > 20 \text{ ps}$), the data fit well to a $t^{1/2}$ dependence, characteristic of transfer of electron kinetic energy to the lattice and diffusion during the pulse. The damage is thermal in nature and characterized by melting and boiling of the surface. For long pulses, heating of the lattice and subsequent thermal damage can occur without significant collisional ionization.13

For real machining applications, the laser intensity is adjusted to be somewhat above the damage threshold. To account for heating of the supercritical plasma, a full radiative hydrodynamic model which accounts for plasma formation, multiple ionization, material equation of state, shock wave generation and cooling by radiation, conduction and plasma expansion (PdV work) is required. A one-dimensional Euler–Lagrangian code was developed for this purpose.26 The predicted pressure and temperature of fused silica machined at the highest irradiance which we have used ($\lambda = 1053 \text{ nm}$, 500 fs, $I = 5 \times 10^{14} \text{ W/cm}^2$) is shown in Fig. 6. The absorbed laser energy heats a skin depth of the material to $\approx 125 \text{ eV}$ at the point of maximum irradiance ($t = 1 \text{ ps}$). There is minimal transport of material during the pulse. As a result, the laser encounters the solid surface for the duration of the pulse depositing energy into solid density material. In the next few picoseconds, this energy is dissipated into the bulk by both shock and electron conduction and is also dissipated by the initial plasma expansion off the surface and radiation. Even under these extreme conditions, a depth of only $\approx 1 \mu \text{m}$ is heated to near 1 eV [Figs. 6(a) and 7(a)]. The edge of the plasma front blows off the surface with a static expansion velocity of $2 \times 10^7 \text{ cm/s}$ over the next nanosecond [Fig 7]. The initial shock wave launched into the material is large reaching 12 Mbar. This shock propagates with a velocity $= 4.8 \times 10^5 \text{ cm/s}$ but dissipates rapidly in magnitude. The pressure associated with the shock drops to less than a few tens of kbar (less than the yield strength of the material) within the first micron [Fig. 7(b)]. This offers extremely high precision machining with no heat or shock effected zone extending beyond $\approx 1 \mu \text{m}$ of the machined surface.

An example of a fused silica surface machined with a single shot is shown in Fig. 8, where the beam intensity was purposefully modulated to demonstrate the threshold nature of material removal in the ultrashort pulse regime. Only in those regions where the intensity is above threshold is any material removed. With all other conditions kept the same but the pulse duration adjusted to 1 ns, the fused silica surface shattered. The nonlinear threshold nature of ultrashort pulse machining enables the production of shaped cuts and holes by shaping the spatial profile of the laser beam. Additionally, high aspect ratio holes ($L/D > 20:1$) can be produced. The aspect ratio and formation of shaped holes are governed by the additional factor of the waveguide nature of the macroscopic dielectric surface. Treatment of these effects is beyond the scope of this article.

The lack of significant energy deposition beyond the volume of interest achieved by using ultrashort pulses enables the use of high repetition (0.1–100 kHz) lasers without the need for external cooling of the part being machined. Even though only a very small depth of material is removed per pulse (0.1–1 \mu m), the high repetition rate enables extremely high cut rates (beyond 1 mm depth per second). The high plasma temperatures achieved within the expanding plasma ensure that the vaporized material will be completely removed from the vicinity of the machined surface before the arrival of the next pulse. For example, an expanding vapor
with even a low expansion velocity of $10^5 \text{ cm/s}$ will be 1 m away from the surface before the arrival of the next pulse 1 ms later (operating at 1 kHz). With conventional nanosecond or microsecond lasers, the vapor will evolve during the laser pulse. For nanosecond lasers, the coupling of the laser light to the solid surface is reduced since the incident laser light will be scattered and absorbed by the vapor plume.

A significant feature of this nonthermal material removal mechanism is that since there is minimal energy deposition in the remaining material, there is a minimal increase in temperature. Thermal measurements show that when irradiated with conventional nanosecond laser pulses, the bulk temperature of a 1 mm slice of tooth increased by over 40 °C.
While for femtosecond pulses the temperature rise was less than 2 °C (Fig. 9). The fluence in each case was set to remove approximately 1 μm depth of material per pulse. This required 30 J/cm² for the ns pulses and only 3 J/cm² for the fs pulses. The practical consequences in dentistry are substantial. In the case of existing laser systems, active cooling of the tooth is necessary to prevent permanent damage to the pulp and nerves which occurs at an increase of ≈5 °C over body temperature. With ultrashort laser pulses, no cooling would be necessary.

Another dramatic example is the machining of explosives with ultrashort laser pulses. Cutting and machining operations on energetic materials present significant safety challenges. With conventional machine tools, improper fixturing of the workpiece, improper tool configuration, and improper cutting speeds have resulted in violent reactions during machining operations. Conventional laser pulses are often used to ignite explosives. However, with femtosecond pulses, plasma formation and material removal occur too fast for significant energy transfer to the remaining material. Furthermore, since only a small amount of material is removed per laser pulse, there is negligible shock imparted to the remaining material. The shock wave that does exist decays to an insignificant level within ≈1 μm of the surface. The waste products from short-pulse laser cutting are, for the most part, solid carbon or benign gases, which can be released into the atmosphere.

We have machined a variety of high explosives including LX-14 (95.5% HMX/4.5% Estane), LX-15 (95% HNS/5% Kel-F), LX-16 (96% PETN/4% FPC 461), LX-17 (92.5% TATB/7.5% Kel-F), PBX-9407 (94% RDX/6% Exxon 461), and pressed TNT. In some of the experiments the beam first cuts through the explosive and then into a stainless steel substrate and in other experiments the beam first cuts through stainless steel and then into the explosive. Figure 10 shows two cuts across a 1-cm-diam, 2-mm-thick pellet made using a 1 kHz, 100 fs Ti:sapphire laser system. Fourier transform infrared spectroscopy of the laser cut LX-16 surface showed no evidence of any chemical reaction products. The laser cut surface was chemically identical to the original LX-16 material. We have also produced very high aspect ratio cuts and holes (~1000:1) in high explosives. When we used pulses only modestly in the conventional regime (~600 ps) deflagration of the LX-16 pellet was immediately observed. Examination of the pellet afterward revealed that the edges of the cut were melted and contained a multitude of reaction products consistent with thermally induced ignition [Fig. 10(b)].

A final example of the ability of intense femtosecond pulses to machine dielectric materials which are extremely difficult to machine by other methods is the class of very hard dielectrics such as diamond, sapphire, and silicon carbide. We have machined all of these materials with ease. A 500-μm-diam hole produced through 2 mm of SiC with 120 fs pulses focused to a peak fluence of ≈3 J/cm² is shown in Fig. 11. In fact, SiC exhibits one of the lowest thresholds and highest machining speeds with ultrashort pulses we have observed in any material. When the pulse duration was increased to even a few tens of picoseconds, the machining threshold increased dramatically and significant thermal shock (cracking) was observed.

**SUMMARY**

In summary for dielectric materials, there is a strong deviation from the usual $t^{1/2}$ scaling of laser damage fluence...
for pulses below 10 ps. The damage threshold continues to decrease with decreasing pulse width, but at a rate slower than $r^{1/2}$ in the range 0.1–10 ps. As the pulses become shorter than $\approx 0.1$ ps, the increasing influence of multiphoton ionization results in a rapid decrease of the damage threshold. Multiphoton ionization provides an upper limit to the damage fluence for short pulses preventing any increase in damage threshold with decreasing pulse duration as might be predicted from a pure avalanche model. This new mechanism of damage (material removal) is accompanied by a qualitative change in the morphology of the interaction site. The damage site is limited to only a small region where the laser intensity is sufficient to produce a plasma with essentially no collateral damage. This process enables high precision machining of all dielectrics with no thermal shock or distortion of the remaining material. Although the absolute machining fluence varies, all pure dielectrics (oxides, fluorides, explosives, plastics, glasses, ceramics, SiC, etc.) exhibit similar behavior and can be machined to micron scale precision.

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26. We have modified and extended the HYADES radiation hydrodynamics code developed originally by J. T. Larsen, and described in *Radiative Properties of Hot Dense Matter* (World Scientific, Singapore, 1991).