

Integrated Laser-Material Interaction and PIC-DSMC Simulation of Laser-Triggered Vacuum Switches

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Abstract—This article describes a numerical investigation of laser-triggered, high-voltage switches. Particle-in-cell Monte Carlo (PIC-DSMC) simulation is combined with quantum-mechanical modeling of light-material interaction to provide a spatiotemporal description of the discharge process during and after a laser pulse impinges on an electrode. The light-material interaction model incorporates a microscopic treatment of laser energy absorption by the electrode material, as well as the subsequent electron kinetics that includes the influence of the electrode electronic band structure and relaxation processes from carrier-carrier and carrier-phonon scattering. This light-material interaction model is one-way coupled into the PIC-DSMC calculations as a precomputed, time-varying, material injection boundary condition, allowing for more accurate simulation of the plasma evolution in the switch as the laser impacts the electrode surface and a subsequent conductive channel is formed between electrodes. Parametric studies of switch trigger time over ranges of laser pulse duration and pulse energy are performed. The results indicate that breakdown times on the order of 10's of nanoseconds are possible using commercially available laser diodes with moderate power densities of $\sim 10^{11}$ W/m².

Index Terms—Cathode vaporization, high-voltage switching, laser-material interaction, laser-triggered vacuum switches (LTVSs), particle-in-cell direct simulation Monte Carlo, plasma breakdown, pulsed power systems, thermionic emission.

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I. INTRODUCTION

INVESTIGATIONS of laser triggering of electrical switches span a long period [1], [2], [3], [4]. High-voltage switches are used in many electrical systems such as pulsed power technologies, electric vehicles, power transmission, and sustainable energy. The need for better performance, especially high current transmission with fast delay time, is more important than ever, as modern society continues to transition from fossil to electrical power sources. For certain switching applications, laser triggering is gaining interest because of its potential to significantly improve operating range, timing reliability, and switch lifetime. In particular, there is a growing desire to realize such switches using commercially available semiconductor laser diodes, which offer advantages in size, robustness, efficiency, and cost compared to high-power laboratory laser systems. However, it remains unclear whether such moderate-energy sources can reliably produce the material emission and plasma conditions required for fast switch triggering [5], [6], [7], [8], [9], [10], [11].

The dynamics of a laser-triggered switch span many physical regimes, from optical absorption in the solid and thermal transport in the electrode to plasma formation in the gap and eventual breakdown. Prior modeling efforts have typically examined these mechanisms in isolation or with reduced-order descriptions. For example, Zhou and Zhang [11] developed a quantum model for photoelectron emission from metals that captures nonequilibrium electron heating under short laser pulses, while [8] analyzed the sensitivity of material ablation to thermal conduction and laser energy. Experimental and modeling studies of laser-triggered vacuum switches (LTVSs) [3], [12], [13] have investigated plasma formation initiated by laser interaction with a separate target material. In these systems, relatively high-energy laser pulses are used to vaporize and ionize the target, producing a plasma that expands across the electrode gap. While these approaches capture important aspects of laser-triggered switching, they do not provide a predictive, self-consistent description of how laser parameters determine the coupled processes of electron emission, material vaporization, plasma formation, and eventual breakdown. Reduced-order models, such as the 0-D discharge model in [10], can reproduce breakdown characteristics but rely on prescribed emission or material source terms. More

detailed plasma simulations, on the other hand, typically assume predefined injection of charged or neutral species and therefore do not explicitly resolve the underlying laser-material interaction physics that generates these sources.

The next challenge is transitioning from laboratory demonstrations using high-energy or short-pulse laser systems to practical devices that are compact, robust, and low-power for real-world applications. Advances in semiconductor technology, in terms of spatial footprint, output power, and reliability [14], [15], [16], make it feasible to consider laser-diode-triggered switches as fieldable components. Addressing this challenge requires a predictive understanding of how laser parameters accessible to diode systems translate into electron emission, material vaporization, and ultimately plasma formation and breakdown [17]. When a laser trigger pulse impinges on an electrode, it creates a light-matter interaction region in which the laser field modifies the electron energy distribution. Free-carrier absorption promotes electrons to higher energy states with phonon assistance, while interband absorption generates electron-hole pairs that may subsequently recombine via spontaneous emission.

The response of the electrode to the laser pulse depends on the relative timescales of several competing processes. Optical transitions by themselves produce nonequilibrium carrier population distributions and, if the laser frequency is high enough, cause direct photoemission over the electrode work function barrier. Additionally, in the light-matter region, there can be laser-driven thermionic emission and field emission (due to switch holdoff voltage). Intraband scatterings tend to return the carrier populations within each energy band to equilibrium and occur among carriers and between carriers and phonons [18], [19]. In a typical laser-driven switch's electrode material, where carrier densities are relatively high, carrier-carrier scattering occurs at a timescale of tens to hundreds of femtoseconds and drives the nonequilibrium carrier distributions created by optical interactions to quasi-equilibrium distributions at elevated temperatures (carrier heating). Then, at a picosecond time scale, carrier-phonon scattering relaxes the carriers to the lattice temperature. The transfer of energy from the electron-hole plasma to the lattice (lattice heating) is via emission of phonons, and sufficiently high laser power densities can result in high carrier and lattice temperatures, yielding thermionic electron emission and electrode material vaporization [20], [21]. Finally, the laser-driven emission of electrons and vaporized electrode material leaving the interaction region initiates the plasma formation in the gap between electrodes, which eventually results in switch breakdown.

Here, we develop a theoretical and numerical model for a laser-triggered switch tailored to the moderate-energy, nanosecond-microsecond laser pulse regime characteristic of semiconductor laser diodes, to understand the interplay between its light-matter interactions, electron relaxation processes, thermal ablation, and the self-consistent plasma formation and expansion across the gap. In particular, we develop a two-stage framework, wherein a quantum mechanical model is first used to describe the behavior of the carriers in the cathode material as it is illuminated by an external laser, followed by self-consistent particle-in-cell Monte Carlo (PIC-DSMC)

simulations of the plasma that forms between the electrodes due to the injection of material from the quantum model as a boundary condition at the surface. This framework directly links laser parameters to plasma formation without relying on prescribed emission or material source terms. In the first stage, we numerically observe electron thermionic emission and the vaporization of the cathode material. Subsequently, the PIC-DSMC simulations reveal that a density of $> 10^{23} \text{ m}^{-3}$ of ejected cathode material in the gap between the electrodes is needed to close the switch. Our results show that the switch trigger time is essentially independent of the incident laser energy above a threshold of 0.1 mJ that is needed to vaporize sufficient cathode material. Additionally, the lower bound on the switch trigger time is set by the time scale needed for the cathode material to fill the gap between the electrodes, such that shorter laser pulses do not continuously decrease the switching time. More broadly, our results demonstrate that commercially available laser diode systems can provide sufficient energy and power to achieve reliable, fast triggering of high-voltage switches, thereby supporting their use in compact and fieldable pulsed-power applications. Looking forward, this work justifies experimental efforts to build a low-energy laser-diode-triggered electrical switch, and our model framework will enable the study and optimization of the laser, materials, and geometry of the switch.

The remainder of this study is organized as follows. Section II-A reviews the microscopic treatment of light-matter interaction and electron kinetics used in the modeling of the processes resulting from the laser pulse. Section II-B describes the incorporation of the microscopic results into a PIC-DSMC code to track the plasma between electrodes, from formation to eventual switch breakdown. Section III describes a parametric study involving a near-infrared (850 nm wavelength) laser pulse incident on a graphite electrode of a low-pressure, high-voltage switch. A parametric study relates the fluxes and energies of photoelectrons and vaporized electrode material to the dependences of switch trigger time on laser pulse energy and duration.

II. MODEL

Overall, our numerical approach consists of two main parts: a treatment of the electron dynamics in the cathode material using their quantum mechanical equations of motion, which then acts as a source for PIC-DSMC simulations of the plasma that is created within the switch. The first stage takes as input a single-particle band structure, which may be generated using density-functional theory (DFT) methods. Here, these single-particle bands are assumed to be constant as the electrons become excited, as the screening of the Coulomb interaction at high energies typically makes renormalization effects negligible [22]. The second stage (PIC-DSMC) then uses the now-modeled nonequilibrium electron behavior and neutral vaporization as a boundary source for electron and neutral emission into the switch gap; electron ionization of the neutrals results in a conductive plasma that grows across the gap and triggers the switch. By relying on a quantum mechanical description of the electrons, we both avoid the need to parameterize a phenomenological model and are

sensitive to nonequilibrium phenomena that lie beyond such models. This two-stage framework is built on the separation of timescales between ultrafast carrier dynamics (ps) and plasma evolution (ns– μ s), allowing sequential coupling. Importantly, both components have been independently validated: the carrier dynamics model against time-resolved experimental measurements, and the PIC-DSMC implementation through prior verification and validation studies of kinetic plasma behavior [23].

It is important to note that the present model is formulated for moderate-power density ($\sim 10^{11}$ W/m²), laser-diode-relevant regimes, where plasma formation is initiated through thermionic electron emission and laser-induced material vaporization from the cathode surface. This mechanism differs fundamentally from that observed in experiments employing ultrafast or high-intensity laser systems, or operating under high-field or preionized conditions, where plasma can form directly in the gap through multiphoton or avalanche ionization, or via explosive emission processes at the cathode [8], [13], [24], [25]. In such cases, breakdown may occur on sub-10–50 ns timescales without requiring neutral material transport across the gap. These regimes typically involve laser intensities several orders of magnitude higher than those accessible to semiconductor laser diodes or rely on preexisting gas or plasma that alters the discharge dynamics. Accordingly, the present model does not attempt to capture these high-field or ultrafast triggering mechanisms, but instead focuses on the diode-relevant regime.

Additionally, in the conceptual switch design explored here, the carbon cathode directly absorbs the laser energy and vaporizes. This is different from prior switches that use laser power densities $> 10^{12}$ W/m² and operate by irradiating a potassium chloride and titanium cathode, where the laser energy is absorbed by the titanium and then the potassium chloride heats up and vaporizes [26]. Models of the KCL-Ti switch design include a multiphysics description of the electrode heating and dynamic evolution as the laser power density is large enough that one may expect nonnegligible contribution from advection of material on the timescale of the laser [27]. In the present proposed switch model, the power densities are low enough that, on the timescale of the laser, we expect cathode material supply to be dominated by the vaporization, and we neglect material advection.

A. Laser-Material Interaction

We approximate the electrode material to have two parabolic bands separated by a bandgap, as shown in Fig. 1. Moreover, four assumptions are made to lessen numerical demands during parametric studies. First, the light-matter interaction and carrier kinetics may be described in terms of properties that are spatially averaged over an interaction region; this region is defined by the laser beam diameter and an effective penetration depth. The second assumption is that fast carrier-carrier scattering maintains equilibrium within each energy band, so that carrier populations may be described by Fermi-Dirac functions, determined by chemical potentials and temperatures extracted from the instantaneous electron and hole total densities, and using conservation rules of

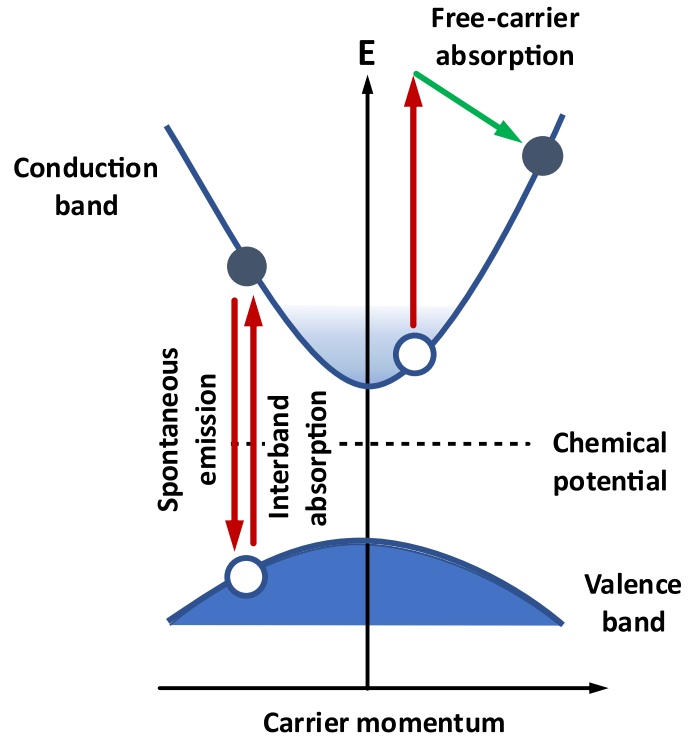


Fig. 1. Schematic of light-material processes and representative electrode band structure. For the graphite cathode treated in Section III, the conduction and valence band dispersions are identical parabolas, and there is no bandgap [31], [32], [33].

intraband scattering. The third assumption is that the effects from the laser pulse (absorption and subsequent interband carrier dynamics) occur on a shorter timescale (0.01–1.0 ps range, [28] [29]) than the plasma evolution between electrodes (1.0–100's ns range depending on gap size, [30]). This decouples the microscopic calculations and PIC-DSMC simulations and allows them to be performed sequentially. Finally, the present model assumes that the spatially varying deformation of the laser switch surface (and the subsequent effects on electron emission and material vaporization) is negligible during the pulse. For relatively uniform laser profiles, this assumption should be reasonable except near the edges. Future work will address this approximation and allow for the surface to deform nonuniformly due to the laser spot profile, accounting for field enhancement and the resultant increase in electron emission, joule heating, and material evaporation. Likewise, there are potentially substantial benefits to be considered in the future that may be realized through the incorporation of patterned surfaces, such as metasurfaces or photonic crystal slabs, to enhance the light-matter interactions in our system.

With the above assumptions, the system's evolution may be tracked by numerically solving the equations of motion for the electron and hole densities (N_e and N_h , respectively), and the laser photon density (N_p) inside the interaction region [34]

$$\frac{dN_e}{dt} = R_{abs} - R_{spont} - R_{emit} \quad (1)$$

$$\frac{dN_h}{dt} = R_{abs} - R_{spont} \quad (2)$$

$$\frac{dN_p}{dt} = P_v - R_{abs} - R_{fca}. \quad (3)$$

Also solved are the total carrier and lattice energies (E_{eh} and E_l , respectively)

$$\frac{dE_{eh}}{dt} = S_{fca} + S_{abs} - S_{spont} - S_{c-p} - S_{emit} \quad (4)$$

$$\frac{dE_l}{dt} = S_{c-p} - S_{cond} - S_{vap}. \quad (5)$$

In the above equations, R denotes a total density change rate, S denotes an energy change rate, and the subscripts label the contributions from interband absorption (*abs*), free-carrier absorption (*fca*), spontaneous emission of radiation (*spont*), electron emission from the material (*emit*), carrier-phonon scattering (*c-p*), thermal conduction from interaction volume to rest of the electrode (*cond*), and vaporization of electrode material (*vap*). For the laser photons, P_v is the rate of photon density introduced into the interaction volume by the laser pulse. Equation (3) indicates that every electron-hole pair creation or free-carrier absorption leads to annihilation of a laser photon.

The derivation of (1)–(5) also gives analytical expressions for the rates (see Appendix A). We use effective relaxation rates to account for the changes in the carrier density distributions caused by carrier scattering [35]. For laser pulses longer than tens of picoseconds, the more rapid carrier-carrier and carrier-phonon scattering relax the carrier distributions to be Fermi-Dirac functions, $f(\varepsilon_{\sigma k}, \mu_{\sigma}, T_c)$ and $f(\varepsilon_{\sigma k}, \mu_{\sigma}^L, T_L)$, which define temperatures for the carrier and lattice, T_c and T_L , respectively. The carrier temperature is calculated from energy conservation in carrier-carrier scattering, i.e.,

$$E_{eh} = \sum_k [(\varepsilon_g + \varepsilon_{ek}) f(\varepsilon_{ek}, \mu_e, T_c) + \varepsilon_{hk} f(\varepsilon_{hk}, \mu_h, T_c)] \quad (6)$$

where ε_g is the bandgap energy, $\varepsilon_{\sigma k}$ is the electron or hole energy in momentum state $\hbar k$, and the summation is over all states. The lattice temperature is determined by (5), using

$$E_l = C_p T_l \quad (7)$$

where C_p is the heat capacity of the interaction volume. For the Fermi-Dirac functions, the chemical potentials are determined by total carrier population conservation in intraband collisions as

$$N_{\sigma} = \frac{1}{V} \sum_k f(\varepsilon_{\sigma k}, \mu_{\sigma}, T) \quad (8)$$

where V is the interaction volume.

There are a few mechanisms that influence the lattice temperature in the interaction region: the transfer of energy from the electrons and holes to the lattice via carrier-phonon scattering, the conduction of heat from the interaction volume to the rest of the cathode, and the energy lost from vaporization of cathode material. These effects, in conjunction with energy conservation, determine the form of (5), where the first term on the right-hand side is from energy infusion via carrier-phonon scattering, the second term is from conduction, and the third from material loss. The expressions for computing S_{c-p} , S_{cond} , and S_{vap} are in Appendix A.

When free-carrier absorption sufficiently raises the carrier temperature T_c , the hottest electrons escape the electrode by

thermionic emission or tunneling. We account for this loss of hot electrons from the cathode (or equivalently, gain of photoelectrons in the gap between electrodes) with

$$\left. \frac{d}{dt} f(\varepsilon_{ek}, \mu_e, T_c) \right|_{emit} = -\gamma_k^t f(\varepsilon_{ek}, \mu_e, T_c) \quad (9)$$

where γ_k^t is an effective electron escape rate (see Appendix A for expression and discussion).

Lastly, supplying vaporized cathode material into the switch gap is crucial for near-vacuum switches to break down and conduct significant current across the switch. The flux and the velocity of the vaporized material affect the neutral density distribution across the gap, which is important for starting the discharge. They are determined from the Hertz-Knudsen equation [21] and S_{vap} in (5). The details are presented in Appendix A. Finally, we note that this carrier dynamics model has been previously verified against experimental results (see Appendix B).

B. PIC-DSMC Simulation

The second half of the model describes the plasma dynamics as the switch closes. As stated earlier, one of the present assumptions is that the laser-material interaction fluxes produced by the laser trigger are independent of the plasma dynamics. This assumes that the plasma does not absorb or otherwise modify the laser pulse and that plasma-electrode interactions (e.g., sputtering, ion-induced secondary emission) either occur after the laser pulse or do not substantially modify the electron or lattice energy distributions while the laser pulse is active. If these conditions are met, then the laser-electrode interaction may be treated as an uncoupled, time-varying, particle injection cathode boundary condition in a PIC-DSMC simulation.

For the present simulations, we use Aleph, a 3-D unstructured finite element ES PIC-DSMC code that is particularly suited for complex device-scale (cm^3) geometries, as it is massively parallel and scales to tens of thousands of CPU cores with dynamic load balancing. PIC-DSMC is a kinetic method for simulating nonequilibrium plasmas [36] with complicated plasma chemistry [37], plasma-surface interaction and emission models, self-consistent plasma-generated photon transport, and dynamic particle reweighting/merging to control particle counts [30], [38], [39]. In Aleph, computational particles each represent a different number of individual physical particles to enable the resolution of the several orders of magnitude differences in densities present in discharges. The computational particles are pushed using a second-order velocity Verlet scheme and a constant E-field in the element; additional details can be found in [39]. To describe the breakdown in the switch gap, these simulations use a relatively simple plasma-neutral interaction model including electron-carbon elastic, excitation, and ionization collision chemistry shown in Table I. Elastic and resonant charge exchange ion-neutral collisions are included using VHS [37] and Smirnov [40] models for the cross sections, respectively.

Aleph has been extensively verified and validated in prior studies of low-pressure, nonequilibrium plasma discharges, providing confidence in its application to the present switch

TABLE I
INCLUDED PARTICLE INTERACTIONS

Interaction	Type	$E_{\text{threshold}}$ (eV)	Reference
$e^- + C \rightarrow e^- + C$	Elastic	-	[46]
$e^- + C \rightarrow e^- + C^*$	Excitation	1.26	[47]
$e^- + C \rightarrow e^- + C^{**}$	Excitation	2.68	[48]
$e^- + C \rightarrow 2e^- + C^+$	Ionization	10.6	[46]
$e^- + C^+ \rightarrow 2e^- + C^{++}$	Ionization	24.4	[47]
$e^- + C^{++} \rightarrow 2e^- + C^{+++}$	Ionization	41.4	[47]
$C + C^+ \rightarrow C^+ + C$	Charge Exchange	-	[40]
$C + C^+ \rightarrow C + C^+$	Elastic	-	[37]
$C + C \rightarrow C + C$	Elastic	-	[37]

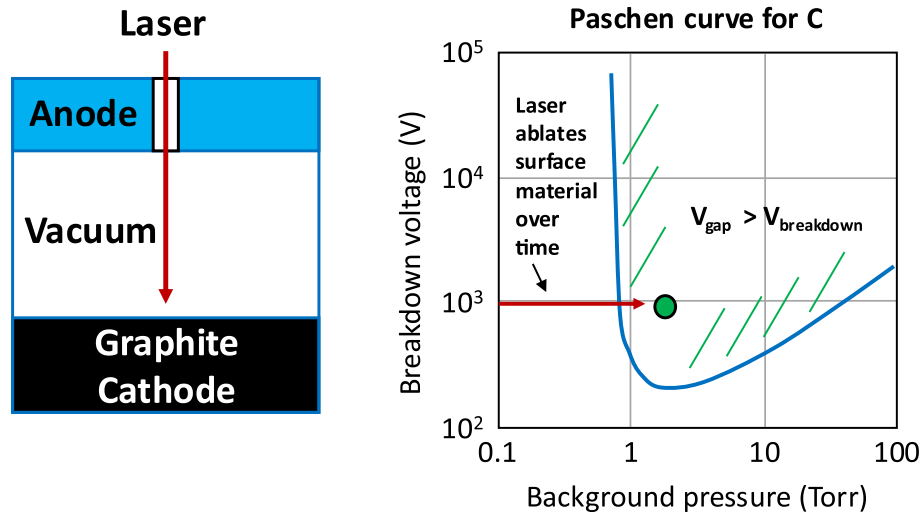


Fig. 2. (Left) Basic features of a laser-triggered switch. In the modeling, we assume a vacuum gap of 1 mm. (Right) Schematic of the switch operation showing the pressure increase due to laser ablation, resulting in the anode voltage exceeding the breakdown voltage.

conditions. In particular, simulations of electron transport and sheath formation driven by thermionic emission have shown quantitative agreement with laser-induced fluorescence measurements of electron density [41]. Additionally, simulations of pulsed discharges in helium reproduce experimentally observed nonthermal emission spectra [42]. These validation cases probe quantities that are highly sensitive to the underlying numerical methods, including collision models, charge-to-mesh coupling, and self-consistent electric field evolution. Further verification and validation studies, including assessments of stochastic algorithm robustness and numerical convergence, have also been reported elsewhere [43]. Taken together, these results establish that Aleph accurately captures the kinetic plasma behavior relevant to the breakdown dynamics modeled in this work.

For self-sustaining breakdown situations, ion- and photon-induced electron emissions from the cathode are taken into account. The present simulations use a constant ion-induced secondary electron emission (SEE) (ion-induced SEE) coefficient, $\gamma_{i-SEE} = 0.1$, a representative value for <100 eV ions hitting metal surfaces [44], and a wavelength-dependent photoemission yield where the photons tracked in the PIC-DSMC simulation are the result of electronically excited neutrals that have spontaneously transitioned to a lower electronic state [38]. Information on radiative lifetimes of excited carbon states is provided using the NIST database [45]. Neutral carbon atoms incident on the electrode surfaces scatter diffusely, based on a 300 K surface temperature for the anode and a time-varying temperature based on the surface-material interaction model's lattice temperature.

III. APPLICATION OF APPROACH TO LASER TRIGGERED SWITCH

To demonstrate the combined laser-material interaction and PIC-DSMC approach, we apply it to model the triggering of a high-voltage vacuum switch. Fig. 2 shows the modeled laser-triggered switch and the operational principle of the switch. In the absence of laser vaporization of cathode material into the gap, the breakdown voltage is ~ 10 kV, and the anode is held at 1 kV. The 850 nm wavelength laser pulse from a semiconductor laser array enters through a $200\text{ }\mu\text{m}$ diameter aperture in the anode. Subsequent vaporization of the graphite cathode increases the background pressure (Fig. 2, right) in the switch gap, resulting in breakdown as the anode voltage eventually exceeds the breakdown voltage. The present work assumes an initial vacuum; however, the model presented here can be applied to a switch with an initial background gas as long as any additional plasma chemistry for the initial background gas is added, provided it is a different species than the electrode material. The assumptions for an initial gas are that the switch's initial pressure times the gap length ($p d_{\text{gap}}$) is on the left side of the Paschen curve (Fig. 2, right), the switch holdoff voltage is below the vacuum breakdown voltage for the given surface roughness, contamination, etc. (fields of $10\text{--}100$'s MV/m or a voltage of $10\text{--}100$ kV for our 1 mm gap, [49]), and the initial background gas is transparent to the laser wavelength.

Three-dimensional simulations are computationally expensive, and therefore, to reduce the computational cost for scoping out the feasibility of a laser diode optical switch, we simulate both the plasma and the laser one dimensionally. This neglects geometric effects such as the electrode shape and the plasma expansion away from the laser spot, as well as assuming the laser spot is spatially uniform with a top-hat profile. It should be noted that the model presented here is not limited to these assumptions as we can extend the PIC-DSMC model to 3-D and use spatially varying laser energy fluxes to determine the local material response boundary condition for the PIC-DSMC simulation assuming that the lateral conduction timescale is slow ($\sim 1\text{ }\mu\text{s}$ for a length scale equal to the laser radius) compared to the laser pulse timescale ($\sim 0.1\text{ }\mu\text{s}$).

A. Simulation of Plasma Evolution From Laser-Triggered Initiation to Switch Breakdown

We start by solving (1)–(5) to obtain the carrier populations and energies within the graphite material after the arrival of the laser pulse. We compute the electron and material fluxes stemming from various incident laser beam energies and pulse durations for use in the PIC-DSMC simulations shown later; for brevity, here we only show the results for a laser pulse energy of $\int_{-\infty}^{\infty} dt \hbar \omega N_p(t) = 0.1\text{ mJ}$ and a pulse duration of 100 ns, giving a peak power of 1 kW. As seen in Fig. 3(a), the laser pulse reaches 10% of peak power after ~ 80 ns, reaches peak power at 150 ns, and decreases back to 10% peak power after ~ 250 ns.

Within the interaction volume, the laser pulse absorption creates electron-hole pairs and hot electrons. Fig. 3(b) shows

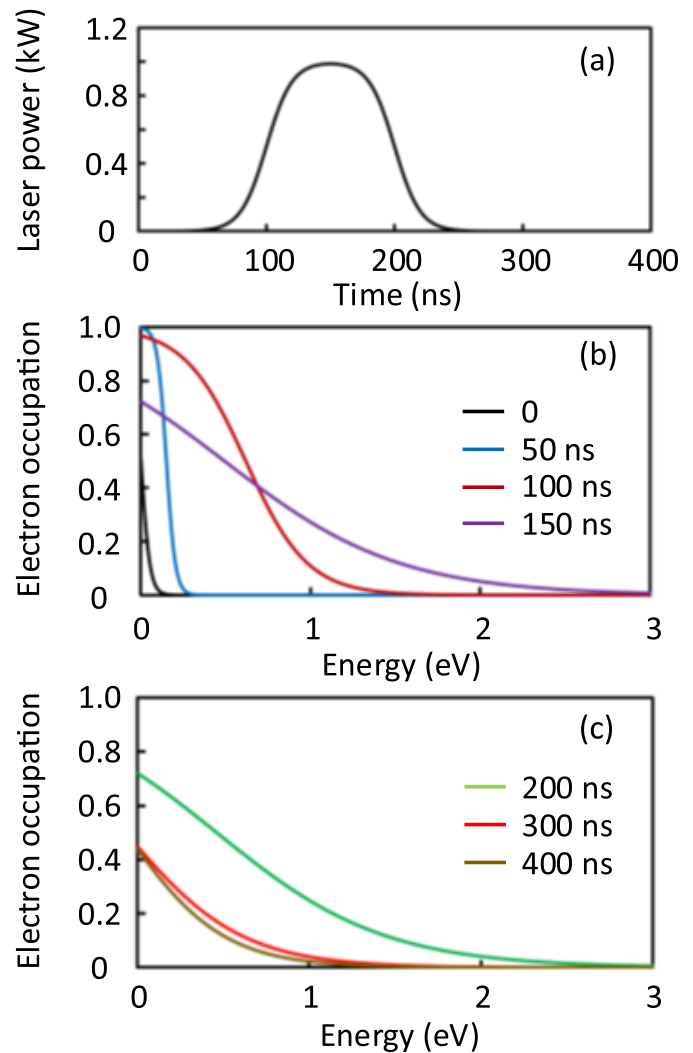


Fig. 3. (a) Laser power as a function of time. Electron distribution (b) before and (c) after the peak of the laser pulse. The times are as labeled, and the electron energy is referenced to the conduction band edge.

the calculated changes in the electron-population energy distribution created by the laser pulse. The black curve is the starting distribution: a Fermi-Dirac function at 300 K with chemical potential at the conduction band edge. At 50 ns, prior to significant laser intensity, changes are already occurring in the electron distribution from the creation of electron-hole pairs (blue curve). From 50 to 100 ns, the electron population increases significantly from interband absorption (compare areas under blue and red curves). Once a hole population is created, recombination of electrons and holes occurs by spontaneous emission in the cathode material. The difference between interband absorption and spontaneous emission contributions gives the net rate of change of electron density. The red curve also indicates filling of lower energy states to the point where Fermi blocking may be significant. After 100 ns, free-carrier absorption starts to dominate, and higher energy states become occupied (carrier heating as shown by the purple curve). Owing to rapid carrier-carrier and carrier-phonon collisions (effective relaxation times of ~ 100 and ~ 1 ps, respectively), the electron distributions at all times

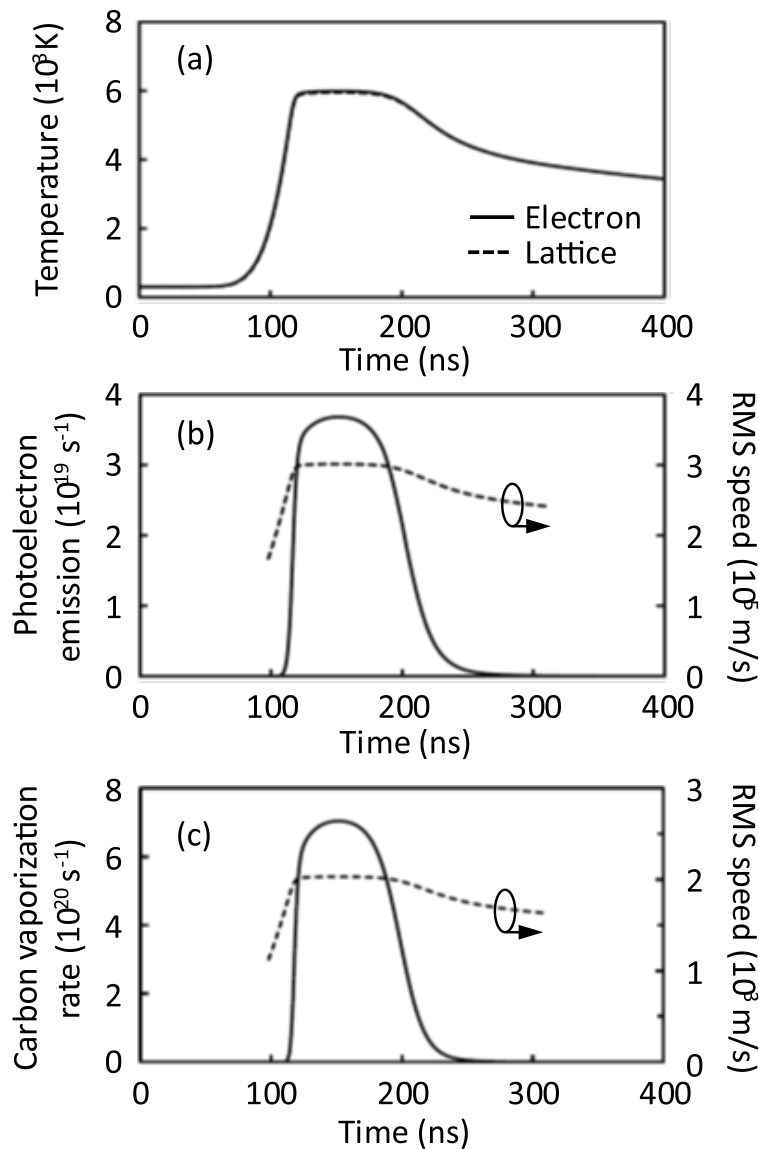


Fig. 4. (a) Electron and lattice temperatures versus time. (b) Rate of electron photoemission (solid curve) and root-mean-square speed of emitted electrons (dashed curve) versus time. (c) Rate of carbon vaporization (solid curve) and root-mean-square speed of emitted carbon atoms (dashed curve) versus time.

are, to a very good approximation, described by Fermi-Dirac functions at different chemical potentials and temperatures.

Between 150 and 200 ns, the system reaches an approximate steady state, where the effects on the electron energy distribution from laser radiation absorption are balanced by photoelectron emission and energy transfer to the lattice via electron-phonon scattering. This is indicated by the similarity of the purple curve in Fig. 3(b) and the green curve in Fig. 3(c). After the passage of the laser pulse, the electron distribution loses population and energy due to spontaneous emission and electron-phonon scattering [red curve in Fig. 3(c)]. Owing to the low intensity of $< 10^7 \text{ Wm}^{-2}$ in a 100 ns 1 mJ laser pulse (for reasonable laser spot sizes on the electrode), electron and lattice temperatures are essentially similar at all times, and they both decrease by cathode material vaporization, electron emission, and heat conduction [brown curve in Fig. 3(c)].

Two microseconds after the passage of the laser pulse, most of the energy remaining in the lattice is transferred to the substrate by heat conduction. For the parameters chosen in this laser-material interaction simulation, this energy transfer occurs around $2 \mu\text{s}$ after the onset of the laser pulse.

The electron dynamics depicted in Fig. 3 drive the production rate and energy of electrons and vaporized carbon entering the vacuum gap between the anode and cathode, via (9) and (A19), and the electron and lattice temperatures in the laser-material interaction region. Fig. 4(a) shows the essentially identical rise in electron and lattice temperatures from the absorption of the laser pulse in Fig. 3 and subsequent electron-electron and electron-phonon scattering. For laser pulse intensities higher than 10^8 Wm^{-2} , there is a noticeable lag in the lattice temperature. The rise in electron temperature leads to thermionic emission and tunneling of electrons into

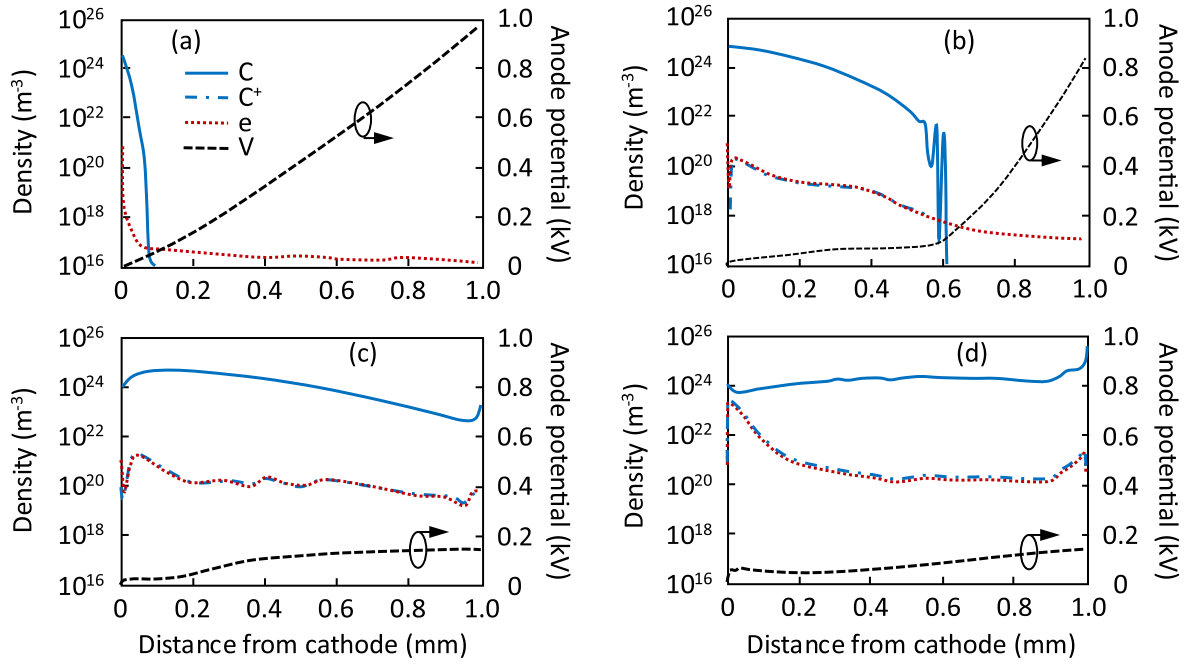


Fig. 5. Carbon atom, carbon ion, and electron densities across the entire 1 mm gap (left y-axis). Also plotted is the gap voltage (right y-axis). The times are (a) 120, (b) 170, (c) 220, and (d) 310 ns, with results averaged over a 0.1 ns window.

the gap, as plotted in Fig. 4(b). The rise in lattice temperature causes vaporization of electrode material into the gap, as plotted in Fig. 4(c).

With the behavior of the electrons and lattice in, and ejected from, the graphite cathode understood, the calculated electron emission and lattice temperature define the boundary conditions for PIC-DSMC simulations that model the formation and dynamics of the plasma in the vacuum gap as the switch breaks down and conducts current. Specifically, our approach uses the emitted electron and carbon atom fluxes, drift velocities, and temperatures computed by the laser-solid interaction model (Fig. 4) from each set of laser pulse parameters simulated. The PIC-DSMC simulations used a mesh size of $0.235 \mu\text{m}$ (to match the Debye length for a plasma temperature of 1 eV and density of $10^{21} \text{ \#}/\text{m}^3$), a timestep size of 10^{-13} s (to resolve plasma frequency for stability and mean collision time and particle-CFL for accuracy), and targeted ~ 100 particles per element, N_p , for each species. The simulation results were found to be insensitive to increasing and decreasing the numerical parameters by a factor of two from their nominal values, as seen in Appendix C. Investigating the sensitivity of the results to material properties will be performed in future work that looks to optimize switch performance.

We assume that the anode is initially held at 1 kV and that the gap is in vacuum before the laser arrives. Furthermore, we consider a 1-D switch geometry with a simple external circuit attached to the anode to limit the current rise and applied voltage to the anode, and therefore apply a finite power to the system during the simulation. Since the area of the electrode is arbitrary in our modeled 1-D system (set to 1 m^2), the gap voltage is simply given by

$$V_{\text{anode}}(t) = V_0 - i_{\text{anode}} R_{\text{circuit}} \quad (10)$$

where the modeled anode current, i_{anode} , and external circuit resistance, R_{circuit} , reduce the applied voltage. The resistance, R_{circuit} , is set at $0.5 \text{ m}\Omega$ to allow the switch current density to reach $\sim 0.1 \text{ MA}/\text{m}^2$ before the anode voltage is significantly reduced. With these parameters, the simulated current density typically grows to $\sim 1 \text{ MA}/\text{m}^2$ during breakdown, which is significantly larger than the space charge limited value of $0.074 \text{ MA}/\text{m}^2$ computed from the Child-Langmuir law [50]. The simulated anode current density, j_{anode} , is smoothed over 0.1 ns to reduce the noise from the number of computational particles used in the simulation.

The results of the PIC-DSMC simulation are plotted in Figs. 5 and 6. Fig. 5 shows the carbon (C and C^+) and electron densities, and voltage across the entire gap at different times. Note that the C^{++} and C^{+++} densities are not shown as they were at least an order of magnitude smaller than C^+ and their dynamics are not significantly different from C^+ ; however, the current analysis did not determine if the switch dynamics were sensitive to the inclusion of double and triple ionization interactions. Starting with Fig. 5(a) for 120 ns, the solid blue curve shows a low number of vaporized carbon atoms concentrated within 0.1 mm from the cathode. Negligible ionization of the carbon has taken place because the neutral carbon density is $< 10^{23} \text{ m}^{-3}$ (ionization mean free path $\sim 1 \text{ mm}$) in the region of the gap where electrons from the cathode have sufficient energy (10.6 eV) for ionization to occur and form a quasi-neutral plasma. The electron flux is space charge limited, and a virtual cathode forms that turns the majority of photoelectrons back into the actual cathode. The electrons that do escape the virtual cathode rapidly accelerate across the gap without collisions because of the low carbon density, except for very near the cathode, where the virtual cathode keeps the electron energies low. Additional results for times earlier than 120 ns are provided in Appendix C.

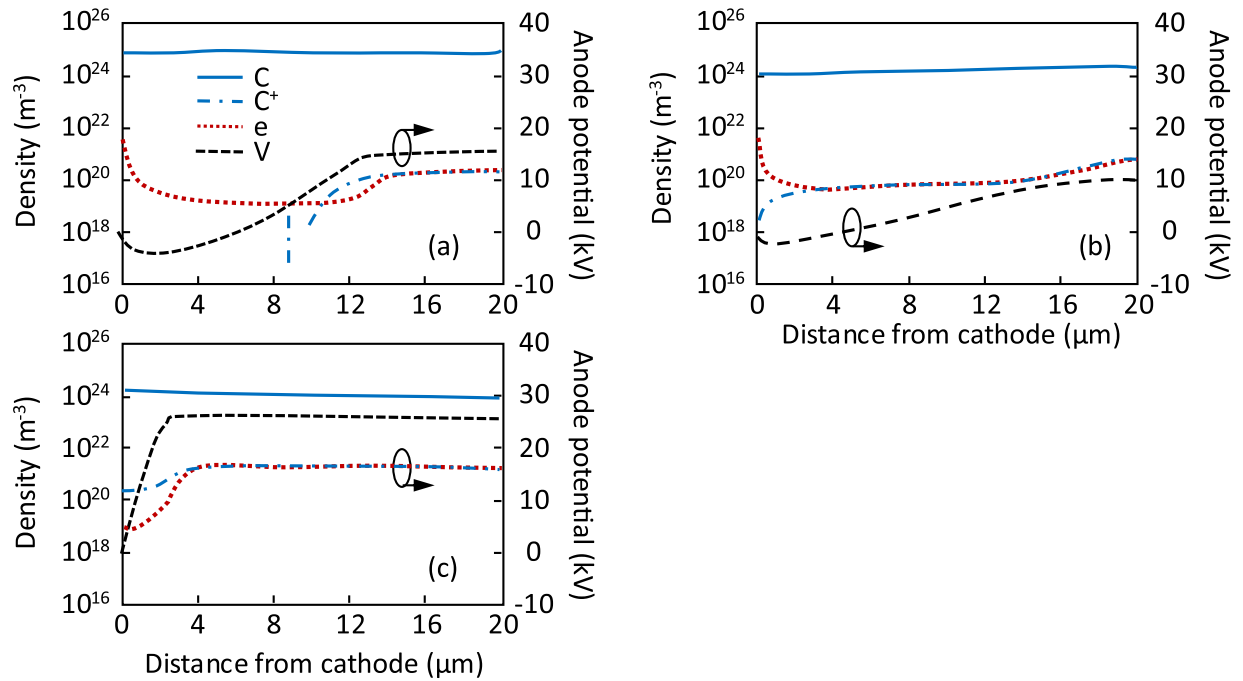


Fig. 6. Carbon atom, carbon ion, and electron densities up to 20 μm from the graphite cathode surface. Also plotted is the voltage. The times are (a) 170, (b) 220, and (c) 270 ns, with results averaged over a 0.1 ns window.

Fig. 5(b) shows the switch's properties for 170 ns after the start of the laser pulse, and reveals increased vaporized carbon density of $> 10^{23} \text{ m}^{-3}$ extending further into the gap. Carbon atom ionization increases appreciably (dot-dashed blue curve), and a quasi-neutral plasma forms up to the edge of the carbon density front as it advances across the gap. The voltage collapses across the quasi-neutral plasma (see dashed black curve at 0.6 mm), and from 120 to 170 ns, there is an evolving plasma sheath that allows higher current (greater than the vacuum space charge limit) to pass through the gap. There is still a virtual cathode extending to 7 μm with a 5 V dip [see black dashed curve in Fig. 6(a)]. This is because photoelectron flux continues to greatly exceed the current that the gap can support. At this time, the carbon ions still do not reach the cathode, as they do not have the energy to get past the virtual cathode established by the excessive injection of photoelectrons [dot-dashed blue curve, Fig. 6(a)]. The carbon ions undergo frequent resonant charge exchange collisions that result in relatively low average ion energies. Therefore, plasma-surface interactions should still not be affecting the laser-material interaction as assumed. In the laser-material interaction model, we also used a constant normal E-field at the electrode surface to compute the contribution to tunneling, and at 170 ns (near the peak of the laser power), we now see substantial changes in the near-cathode E-field. In fact, up to this time, the virtual cathode reversed field would have acted to prevent tunneling. However, based on the low contribution of tunneling current in the laser-surface interaction model, the net result of coupling a time-varying E-field into the laser-surface interaction model would be that the net current into

the gap past the virtual cathode would barely change (instead, we would be over-injecting slightly less).

At 220 ns, neutral carbon fills the entire gap, though not uniformly, as there is increased density closer to the cathode. The plasma density has grown [dot-dashed blue and dotted red curves in Fig. 5(c)] and the virtual cathode, while still present, has shrunk considerably [compare black dashed curves in Fig. 6(a) and (b)]. The carbon transit time of ~ 100 ns across the 1 mm gap is consistent with the laser-material model predicted rms speed of 2000 m/s at a temperature ~ 5000 K. At this stage, the neutral carbon vaporized by the laser reduces, with its density decreasing with rms velocity away from the cathode. Near the anode, the carbon density increases as fast carbon atoms diffusely scatter off the 300 K anode wall and thermalize.

Even here at the tail end of the laser pulse [220 ns, Fig. 5(c)], the photoelectron flux remains space charge limited; however, the current is higher with the quasi-neutral plasma filling the gap and creating a plasma sheath. The small (now ~ -2 V) potential dip continues to prevent a sizable number of carbon ions from reaching the cathode. This is clearly indicated by the dip in ion density (dot-dashed blue) curve at less than 4 μm in Fig. 6(b), where carbon ion density is roughly two orders of magnitude smaller than everywhere else in the gap. The result is that after 220 ns, there is still negligible ion-induced SEE from the cathode. Further into the pulse at 270 ns, the photoemission electron flux falls sufficiently [Fig. 4(b)] causing the virtual cathode to collapse [dashed black curve, Fig. 6(c)]. An ion-rich sheath forms, whose sheath potential is monotonically increasing, notwithstanding small

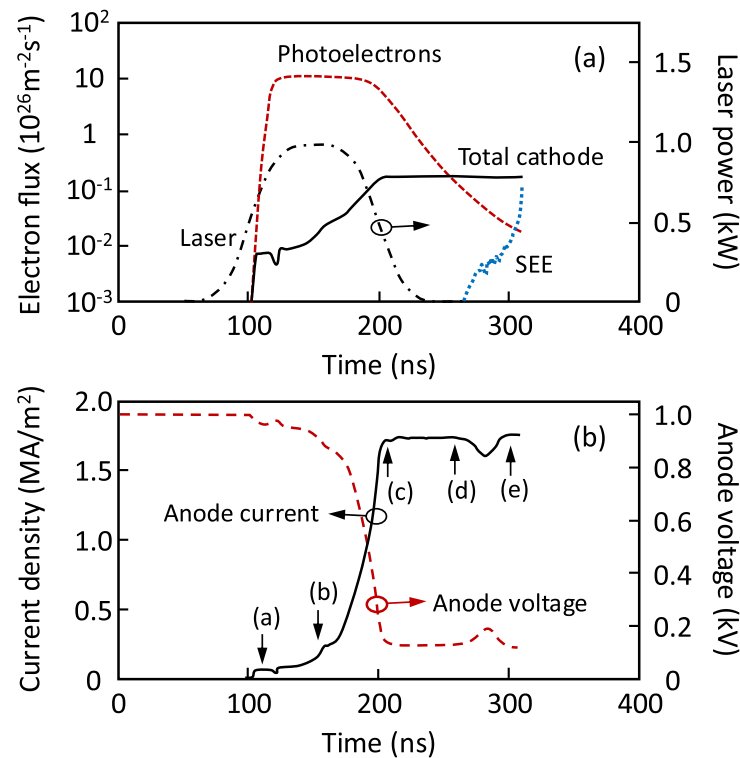


Fig. 7. (a) Trigger-laser pulse and net cathode electron flux (dot-dashed and solid black curves) versus time. Also plotted are the photoelectron flux created from the laser pulse and ion-induced SEE flux (dashed red and dotted blue curves, respectively). Comparison of solid black and dashed red curves indicates that only a small fraction of photoelectrons contributes to the net electron flux in the gap. (b) Switch current and anode potential (solid black and dashed red curves) versus time. The labels [a]–[e] denote the times (120, 170, 220, 270, and 310 ns) for comparison to the plots in Figs. 5 and 6.

noise fluctuations. The ions now reach the cathode surface, leading to ion-induced SEE [see Fig. 7(a)]. Lastly, Fig. 5(d) shows the plasma at 310 ns, where ion-induced SEE dominates the photoelectron flux. As such, the switch discharge becomes self-sustaining without the need for more photoelectrons from the laser pulse.

The details provided in the analyses leading to the results shown in Figs. 5 and 6 allow us to isolate and understand the physical mechanisms that control breakdown evolution in a laser-triggered switch. The black dot-dashed and solid curves in Fig. 7(a) are plots of the time dependences of the laser pulse power and net electron flux emitted by the cathode. The net cathode electron flux shows an initial space charge limited plateau after the laser turns on, and then a gradual rise of the electron flux from 100 to 200 ns as the plasma sheath shrinks during the passage of the laser pulse and the filling of the gap with carbon gas. At 200 ns, the switch has closed, and the anode current plateaus, and the external resistor decreases the anode voltage [solid and dotted curves in Fig. 7(b)] and limits further current growth. For the considered example of 0.1 mJ laser pulse energy, the complete switch closure occurs ~50 ns past the peak of the laser pulse and ~100 ns after significant carbon vaporization flux begins and the carbon crosses the gap.

Contributing to the net electron flux are photoelectrons and ion-induced SEE. The red dashed curve in Fig. 7(a) shows that only a fraction of the photoelectrons contribute to the net

electron flux. This is because, as photoelectron flux rapidly increases past the (initial) 0.074 MA/m^2 space charge limit for the gap, a sizable fraction of the photoelectrons return to the cathode. Between 120 and 200 ns, the cathode material is vaporized into the gap, and a quasi-neutral plasma builds up. The plasma creates a cathode sheath that causes the anode current to grow rapidly for ~50 ns, while the external circuit resistor reduces the anode voltage, as shown in Fig. 7(b). Beyond 220 ns, the photoelectron contribution decreases with the passage of the laser pulse. As the photoelectron current decreases below the space charge limit, carbon ions in the plasma begin to reach the cathode, and ion-induced [SEE, dotted blue curve in Fig. 7(a)] increases to keep the total cathode current constant.

Note that the anode current and voltage remain fairly constant until ~250 ns, despite the termination of laser power and subsequent decrease in the photoelectron flux. Around 250 ns, the anode current starts to dip and then eventually recovers after another ~50 ns. This current dip marks the time when the photoelectron flux drops below the space charge limit, and thus the total current across the gap decreases until sufficient ion flux reaches the cathode, such that ion-induced SEE replaces the dropping photoelectron flux. If we artificially turn off ion-neutral collisions and allow the ions to free-stream back to the cathode, the current dip is reduced. The recovery time is thus determined by both the ion transit time to reach the cathode surface and how fast the photoelectron flux is

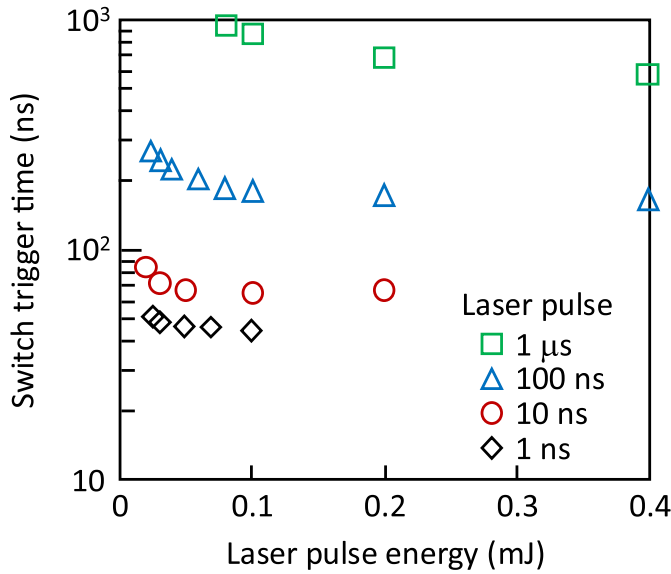


Fig. 8. Switch trigger time versus laser pulse energy from the combined laser-material interaction and PIC-DSMC simulations. The different symbols for laser pulse durations are indicated.

decreasing. Based on the results of Figs. 6 and 7, the timescale for filling the vacuum gap with vaporized carbon is clearly important for switch performance, and thus the drift speed (i.e., mobility) of the vaporized electrode material is a key parameter for future designing optimal switch performance.

After 305 ns, ion-induced SEE from the cathode dominates photoelectron emission, and the switch operates in a more typical Townsend discharge regime, which allows the switch to continue in its conductive state without the laser pulse. In Townsend discharge with a simple external circuit, the voltage across the gap is reduced sufficiently such that the ion-induced SEE rate balances the ionization rate across the gap, i.e., cathode electrons produced per incident ion equals ions produced per cathode electron. With a basic external circuit, this results in a steady-state switch current that is dependent on SEE rate and electron-neutral collision cross sections.

B. Parametric Study of Laser-Triggered Switch Trigger Time

In the second part of this section, we show the application of the two-step laser-material and PIC-DSMC simulation in a parametric study to inform laser triggering and engineering design. Repeating the simulation steps described in Section III-A for a range of different diode laser pulse durations and energies, we evaluate the feasibility of a commercially available laser diode to create the necessary emitted material to produce desirable switch trigger (or function) times. We consider an external circuit where the switch anode current density typically reaches a quasi-steady state of $\sim 1\text{--}2\text{ MA/m}^2$, and define the switch trigger time as the difference between the onset (10% peak power) of the laser trigger pulse and when the anode current density reaches 0.74 MA/m^2 . This current density is ten times the space charge limited current density of the empty gap and roughly half the typical quasi-steady state value.

Fig. 8 plots the results where we vary the laser pulse energy for different pulse durations. The curves indicate that the switch trigger time is relatively insensitive to laser pulse energy for pulse energies above 0.1 mJ, which is the approximate energy needed to vaporize a sufficient amount of electrode material to support self-sustaining Townsend breakdown in the 1 mm gap held at 1 kV. Once there is sufficient laser energy to vaporize the required electrode material, vaporization of additional material does not significantly reduce the time required for switch breakdown. Assuming a simple Paschen-curve breakdown [15] and particle interactions listed in Table I, the gap has to reach a uniform background gas density of $\sim 10^{23}\text{ m}^{-3}$. For a 1 mm gap and assuming a column of vaporized carbon with a radius equal to the 100 μm laser beam radius, the density requirement means that we need to vaporize $\sim 10^{12}$ carbon atoms, equivalent to $\sim 10^{-3}\text{ mJ}$ of deposited energy based on the graphite enthalpy of vaporization. This energy requirement is an order of magnitude smaller than indicated in Fig. 8; the discrepancy stems from the significant laser energy lost to thermal electron emission, kinetic energy of the vaporized carbon, and lattice heating.

Note that the threshold energy depends on laser pulsewidth. With the long microsecond pulse, appreciable energy is conducted into the bulk material, thus reducing the efficiency of vaporizing the electrode material. Consequently, higher laser energy is required for breakdown. For a 1 μs laser pulse (square green points in Fig. 8), there is no switch breakdown below a pulse energy of 0.05 mJ.

Based on the simulations, we find that for most high-voltage switch applications, the laser pulse may be produced by a laser diode array operating in quasi-continuous wave (cw) with direct current modulation. [51] By tailoring the injection current, a commercially available laser-diode bar, capable of delivering 100 W cw out of a 100 μm diameter fiber bundle, may be fine-tuned to produce switch trigger times ranging from tens of nanoseconds to microseconds. With switches operating at a duty cycle of <1 discharge per millisecond, active cooling is not needed to produce laser peak power up to $\sim 1\text{ kW}$ out of the fiber.

Table II summarizes the results where we fix the laser pulse energy to 0.1 mJ, and vary the laser pulse duration from 1 ns to 1 μs . For long ($>100\text{ ns}$) pulses, the switch trigger time nearly equals the laser pulse duration. For instance, the ratio of trigger time to pulse duration is 0.9 for a 1 μs laser pulse. On the other hand, for shorter ($<10\text{ ns}$) laser pulses, the switch trigger time is considerably longer than the pulse duration, i.e., the switch triggers long after the laser pulse is over. With the short laser pulses, the speed at which the switch breakdown occurs is limited by the vaporized material's ability to fill the vacuum gap. Thus, our results predict that further decreasing the incident laser pulsewidth will not substantially improve the switch triggering time. In particular, highly nonuniform gas densities near the cathode inhibit breakdown, and therefore, much of the vaporized cathode material must expand and fill the gap before breakdown can occur. If the cathode material expands away from the cathode at 25 km/s, it will take 40 ns to cross a 1 mm gap. In Table II, this is the reason

TABLE II
TRIGGER TIME VERSUS LASER PULSE DURATION FOR 0.1 MJ LASER ENERGY

Pulse Width (ns)	Trigger Time (ns)	Trigger Time Pulse Width
1	44.7	44.7
10	66.2	6.62
100	186	1.86
1000	900	0.9

for the ratio of trigger time to pulse duration of 44.7 for the 1 ns laser pulse. On the other hand, long pulses allow the vaporized material to equilibrate into a uniform density across the gap on the timescale of the pulse, and thus the switch time is limited by the time it takes for the laser to vaporize sufficient material to support breakdown. As such, we expect pulses longer than 1 μ s to also exhibit switching times similar to the pulsewidth.

IV. CONCLUSION

This article describes a two-stage approach to modeling laser-triggered electrical switches. First, a quantum mechanical model is used to describe interband and intraband absorption of laser radiation by the cathode material, where the absorption, together with carrier-carrier and carrier-phonon scattering, causes electron and lattice heating. The rise in electron and lattice temperatures leads to electron thermionic emission and tunneling, and cathode material vaporization into the gap between electrodes. In the second stage, a PIC-DSMC simulation is used to self-consistently describe the spatiotemporal formation and evolution of the plasma within the gap that eventually results in the closing of the switch. In the present work, we implemented one-way coupling of the laser-material interaction model outputs (fluxes and energies of photoelectrons and vaporized cathode materials) into the PIC-DSMC simulations as injection boundary conditions at the cathode. Compared to prior modeling approaches that rely on prescribed emission or reduced-order descriptions, the present framework enables a direct, physics-based connection between laser parameters, material response, and plasma evolution without relying on empirically prescribed emission or plasma source terms. This capability is particularly important in the moderate-energy regime considered here, where the feasibility of laser-diode-triggered switching depends sensitively on the efficiency of electron emission and material vaporization.

To demonstrate the approach, we analyzed the switch trigger time of a high-voltage, vacuum switch, triggered by an 850 nm wavelength laser pulse incident onto a graphite cathode. Modeling was performed for a range of laser pulse durations and peak powers. The switch trigger time was found to be relatively independent of the laser energy above 0.1 mJ, which is determined by the laser energy needed to fill the gap with a density of $> 10^{23} \text{ m}^{-3}$ of cathode material. In

addition, we found that for short pulses, the switch trigger time is limited by the time it takes for the vaporized carbon to travel from cathode to anode. Our parametric study indicates breakdown times as fast as 10's of ns are possible with commercially available laser diode arrays. In addition, our results reveal that laser-triggered switching can be achieved using commercially available pump lasers.

Future work will focus on extending the present framework to address several important limitations and opportunities. First, the current model employs a one-way coupling between the laser-material interaction and the plasma dynamics. An important next step is the development of a fully self-consistent, two-way coupled model, in which the evolving plasma sheath and associated electric fields at the cathode surface feedback on the emission processes in real time. Such coupling is expected to be particularly important in higher-power or faster-switching regimes, where field enhancement at the surface may significantly modify electron emission and enable nonlinear growth leading to shorter conduction times [52]. Second, the model produces a number of experimentally testable predictions, including the spatiotemporal evolution of neutral and ion densities in the gap and the saturation of trigger time with increasing laser pulse energy. These predictions suggest clear pathways for experimental validation, for example, through laser-induced fluorescence measurements of species densities or interferometric diagnostics of plasma evolution [41], [42], [53]. Such comparisons would enable quantitative assessment of the model and help refine its predictive capability. Finally, the present framework provides a platform for exploring approaches to enhance laser-material coupling in diode-driven systems. In particular, the incorporation of structured surfaces such as photonic crystals or metasurfaces [54] may enable local field enhancement (e.g., via Purcell effects), increasing emission and vaporization rates without requiring higher input laser energy. This represents a promising direction for improving switch performance while maintaining the compact, low-power advantages of semiconductor laser sources.

APPENDIX A

In this Appendix, we present the expressions for calculating the rates appearing in (1)–(5). The derivations may be found in several textbooks [17], [18], [55].

For the interband absorption rate, semiclassical theory gives [55]

$$R_{abs} = \frac{dN_p}{dt} \frac{\mu^2 \nu}{\hbar \gamma A \epsilon_0 c n_b} \times \sum_k [1 - f(\epsilon_{ek} \mu_e T_c) - f(\epsilon_{hk} \mu_h T_c)] \mathcal{L}_\gamma(\nu - \omega_k). \quad (A1)$$

The corresponding energy contribution is

$$S_{abs} = \frac{dN_p}{dt} \frac{\mu^2 \nu}{\hbar \gamma A \epsilon_0 c n_b} \times \sum_k \hbar \omega_k [1 - f_{ek}(\epsilon_{ek} \mu_e T_c) - f_{hk}(\epsilon_{hk} \mu_h T_c)] \mathcal{L}_\gamma(\nu - \omega_k). \quad (A2)$$

In both equations, c and ϵ_0 are the speed of light and permittivity in vacuum, n_b is the background refractive index, A is the illuminated area, ν is the lasing frequency, and μ is the interband transition dipole matrix element. In the summation over carrier states, $\mathcal{L}_\gamma(x) = 1/[1 + (x/\gamma)^2]$ is the Lorentzian line shape function, $\hbar \omega_k = \epsilon_g + \epsilon_{ek} + \epsilon_{hk}$ is the transition energy, and γ is the dephasing rate due to collisions. Inside the square brackets is the population difference between the lower and upper states involved with the laser transition, which depends on the optical transition, scattering, and electron photoemission. Of these contributions, carrier-carrier scattering, occurring at a femtosecond time scale, is by far the fastest. As a result, the carrier distributions relax instantaneously to Fermi-Dirac functions $f(\epsilon_{\sigma k}, \mu_\sigma, T_c)$ with chemical potential and temperature determined by the values of N_e and E_{eh} as they evolve according to (1)–(5).

For the spontaneous emission, a fully quantized derivation gives the rates [55]

$$R_{spont} = \frac{\mu^2 n_b^3}{2\hbar \pi \epsilon_0 c^3} \sum_k \omega_k^3 f(\epsilon_{ek} \mu_e T_c) f(\epsilon_{hk} \mu_h T_c) \quad (A3)$$

$$S_{spont} = \frac{\mu^2 n_b^3}{2\pi \epsilon_0 c^3} \sum_k \omega_k^4 f(\epsilon_{ek} \mu_e T_e) f(\epsilon_{hk} \mu_h T_c). \quad (A4)$$

In the case of free-carrier absorption, [56] we use an expression derived by assuming an effective interaction Hamiltonian for the combination of a direct (in momentum space) optical transition together with an emission or absorption of a phonon to conserve momentum. In this treatment, the rates are given by [34]

$$R_{fca} = \sum_k (R_k^{fca,in} - R_k^{fca,out}) \quad (A5)$$

$$S_{fca} = \sum_k (S_k^{fca,in} - S_k^{fca,out}) \quad (A6)$$

where the in-scattering contributions are

$$R_k^{fca,in} = \frac{dN_p}{dt} \frac{\nu}{\hbar \gamma \epsilon_0 c n_b A} \sum_{k' < k} (\mu_{k,k'}^{fca})^2 \times [f(\epsilon_{ek'}, \mu_e, T_e) - f(\epsilon_{ek}, \mu_e, T_e)] \mathcal{L}_\gamma(\omega_{k,k'} - \nu) \quad (A7)$$

$$S_k^{fca,in} = \frac{dN_p}{dt} \frac{\epsilon_{ek} \nu}{\hbar \gamma \epsilon_0 c n_b A} \sum_{k' < k} (\mu_{k,k'}^{fca})^2 \times [f(\epsilon_{ek'}, \mu_e, T_e) - f(\epsilon_{ek}, \mu_e, T_e)] \mathcal{L}_\gamma(\omega_{k,k'} - \nu). \quad (A8)$$

The out-scattering contributions are

$$R_k^{fca,out} = \frac{dN_p}{dt} \frac{\nu}{\hbar \gamma \epsilon_0 c n_b A} \sum_{k' > k} (\mu_{k',k}^{fca})^2 \times [f(\epsilon_{ek}, \mu_e, T_e) - f(\epsilon_{ek'}, \mu_e, T_e)] \mathcal{L}_\gamma(\omega_{k',k} - \nu) \quad (A9)$$

$$S_k^{fca,out} = \frac{dN_p}{dt} \frac{\nu}{\hbar \gamma \epsilon_0 c n_b A} \sum_{k' > k} \epsilon_{ek} (\mu_{k',k}^{fca})^2 \times [f(\epsilon_{ek}, \mu_e, T_e) - f(\epsilon_{ek'}, \mu_e, T_e)] \mathcal{L}_\gamma(\omega_{k',k} - \nu). \quad (A10)$$

The effective dipole matrix element is

$$\mu_{k',k}^{fca} = \frac{V}{\pi \hbar^2} (2m_c)^{3/2} \sqrt{\epsilon_k + \hbar \nu} G_q \sqrt{n_q} \mu \quad (A11)$$

where V is the interaction volume, m_c is the electron effective mass, and n_q is the phonon population. The Fröhlich coefficient for the interaction of electrons with longitudinal optical (LO) phonons is [18]

$$\hbar G_q = \sqrt{\frac{\hbar \omega_{LO} V_{sq}}{2} \left(\frac{1}{\epsilon_\infty} - \frac{1}{\epsilon_0} \right)} \quad (A12)$$

where ω_{LO} is the LO-phonon frequency, V_{sq} is the Fourier transform of the screened Coulomb potential energy, ϵ_0 and ϵ_∞ are the low- and high-frequency background dielectric constants.

At high electron temperature, electrons escape from the cathode via thermionic emission and tunneling. Guided by the Fowler–Nordheim theory [20] that describes the tunneling, we introduce an effective electron loss rate to approximate the combined effects of tunneling and thermionic emission through a square barrier for simplicity

$$\gamma_k^t = \sqrt{\frac{e V_{bias}}{2 m_c d_{gap}^2}} \left[1 + \frac{W^2 \sinh^2(\zeta d')}{4 \epsilon_{ek} (W - \epsilon_{ek})} \right]^{-1} \quad (A13)$$

where V_{bias} is the bias voltage across a gap of thickness d_{gap} , and W is the work function and $\zeta = (2m_c(W - \epsilon_{ek}))^{1/2}/\hbar$. We introduce a fitting parameter $d' = \xi d_{gap}$ to approximately reproduce the cathode current–voltage curve without the laser trigger from an applied anode potential, resulting in a triangular barrier to tunnel through at the cathode. A value of $\xi = 10^{-3}$ gave reasonable agreement for the range of applied fields we expected ($\sim$ MV/m). Future work will implement a more accurate general thermo-field electron emission model (e.g., Murphy–Good) [57] through a triangular barrier and allow for accurate emission current for any temperature and field without the fitting parameter. This will be more important if we two-way couple the model and have plasma sheaths form, which significantly change the field at the cathode. In terms of γ_k^t , the rates of electron density and energy decrease due to photoelectrons leaving the cathode are

$$R_{emit} = \frac{1}{V} \sum_k \gamma_k^t f(\epsilon_{ek}, \mu_e, T_c) \quad (A14)$$

TABLE III
LASER-MATERIAL INTERACTION MODEL INPUT PARAMETERS

	Value
Photon energy	$h\nu = 1.46 \text{ meV}$
Interaction volume	$d = 100 \text{ }\mu\text{m}; h = 0.1 \text{ }\mu\text{m}$
Graphite specific heat capacity, C_p	0.72 J/gm/C
Graphite density, ρ	2.1 gm/cm^3
Graphite boiling temperature, T_b	4500 K
Graphite boiling pressure, p_b	1 atm
Graphite enthalpy of vaporization, L_v	720 kJ/mole
Frohlich electron-LO phonon coupling	$\hbar G_q = 6.3 \text{ meV}$
LO phonon energy	$\hbar\omega_{\text{LO}} = 36.4 \text{ meV}$
Electron effective mass	$m_e = 0.4m_0$
Optical dipole matrix element	$\mu = e \times 0.2\text{nm}$
Carrier-carrier scattering rate	$\gamma_{c-c} = 10^{13} \text{ s}^{-1}$
Carrier-phonon scattering rate	$\gamma_{c-p} = 10^{12} \text{ s}^{-1}$
Heat conduction rate	$\gamma_C = 2 \times 10^{11} \text{ s}^{-1}$
Ratio of tunneling depth to gap thickness (anchored by experimental switch)	$\xi = 10^{-3}$
Fraction of vaporized target atoms returning to the surface	$\beta = 0.18$

$$S_{\text{emit}} = \sum_k \gamma_k' (\varepsilon_g + \varepsilon_{ek}) f(\varepsilon_{ek}, \mu_e, T_c). \quad (\text{A15})$$

For the processes involving phonons, we assume an effective carrier-phonon scattering rate γ_{c-p} that may be extracted from ultrafast spectroscopy or quantum kinetic calculations. [19] [58] With this approximation, the energy transfer rate from carriers to the lattice is calculated using

$$S_{c-p} = \gamma_{c-p} \sum_{\sigma=e,h} \sum_k \varepsilon_{\sigma k} [f(\varepsilon_{\sigma k}, \mu_{\sigma}, T_c) - f(\varepsilon_{\sigma k}, \mu_l^{\sigma}, T_l)] \quad (\text{A16})$$

where the lattice temperature comes from numerically solving (5)

$$\frac{dE_l}{dt} = C_p \frac{dT_l}{dt} = S_{c-p} - S_{\text{cond}} - S_{\text{vap}} \quad (\text{A17})$$

and C_p is the heat capacity of the interaction volume. For the rate of heat conduction from the interaction volume to the rest of the cathode, we use

$$S_{\text{cond}} = \gamma_c \sum_{\sigma=e,h} \sum_k \varepsilon_{\sigma k} [f(\varepsilon_{\sigma k}, \mu_l^{\sigma}, T_l) - f(\varepsilon_{\sigma k}, \mu_0^{\sigma}, T_0)] \quad (\text{A18})$$

where γ_c is an effective heat conduction rate and T_0 is the heat sink temperature. Lastly, the Hertz-Knudsen equation gives the density flux of the vaporized cathode material [21]

$$\frac{dN_T}{dt} = (1 - \beta) \frac{p_b A}{m_c} \sqrt{\frac{m_c}{2\pi k_B T_l}} \exp\left[\frac{L_v}{k_B} \left(\frac{1}{T_b} - \frac{1}{T_l}\right)\right] \quad (\text{A19})$$

where L_v is the enthalpy of vaporization, β is the fraction of vaporized target atoms returning to the surface, and T_b is the boiling temperature under pressure p_b . To estimate β , we assumed $\beta = 0$ (reasonable for expansion into a vacuum), computed the electron and neutral particle fluxes for the PIC-DSMC boundary condition, and then counted the number of C atoms that returned to the surface in the PIC-DSMC simulation using the computed fluxes. This gave $\beta = 0.18$, which we used for the simulations shown in the present work. The corresponding energy loss rate is

$$S_{\text{vap}} = L_v \frac{dN_T}{dt}. \quad (\text{A20})$$

Lastly, Table III gives the values of parameters in the expressions for the rates. In particular, the height is determined by the penetration depth of the light source. The diameter is

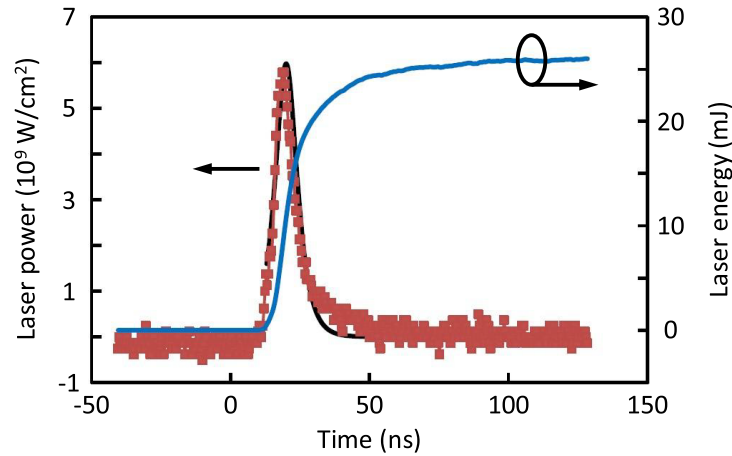


Fig. 9. Laser pulse used in experiment (dots) and modeling (black curve). The pulse energy is 25 mJ (blue curve, right axis) in 10 ns with a beam diameter $220 \mu\text{m}$.

determined through a competition between needing large field intensities, driving toward focusing on smaller spot sizes, and needing to energize a sufficiently large area to eject sufficient material to yield discharge and activate the switch.

APPENDIX B

This Appendix provides validation of the laser–material interaction model against experimental measurements. The goal is to demonstrate that the underlying carrier dynamics, energy transfer, and emission processes used in the main text reproduce experimentally observed behavior in comparable regimes.

An important assumption in the light–matter derivation resulting in (1)–(3) is the quasi-equilibrium approximation, where carrier populations are describable by Fermi–Dirac distributions. This is a reasonable assumption because electron–electron and electron–photon scattering rates ($> 10^{12} \text{ s}^{-1}$) are over three orders of magnitude faster than changes in the electrons and vaporized electrode material, which are governed by the pulse durations and intensities considered in our simulations. Moreover, this assumption has been found valid in experimental data from high-power laser pulse experiments. There, the laser is a solid-state laser, producing pulse durations and intensities considerably shorter and higher, respectively, than laser diodes as are considered in the main text. The theory/experiment comparison also helps indicate the limits of the present treatment.

In the solid-state laser experiment, the pulse energy is $\int_{-\infty}^{\infty} dt \hbar \nu P_V(t) = 25 \text{ mJ}$, duration (full-width half-maximum) is 10 ns, and the spot size diameter is $d = 220 \mu\text{m}$ (Fig. 9). The light is focused onto an Aluminum target. The intensity in the interaction volume is approximately 2.5 MW, which is over three orders of magnitude above the interaction volume considered in the simulations shown in Fig. 3(a).

Numerically solving (1)–(5) for the laser pulse produces the time evolution of the electron and lattice temperatures (inside the interaction volume), the ionized electron population from thermionic emission, and the vaporized Aluminum population (Fig. 10). Input parameters are summarized in Table IV.

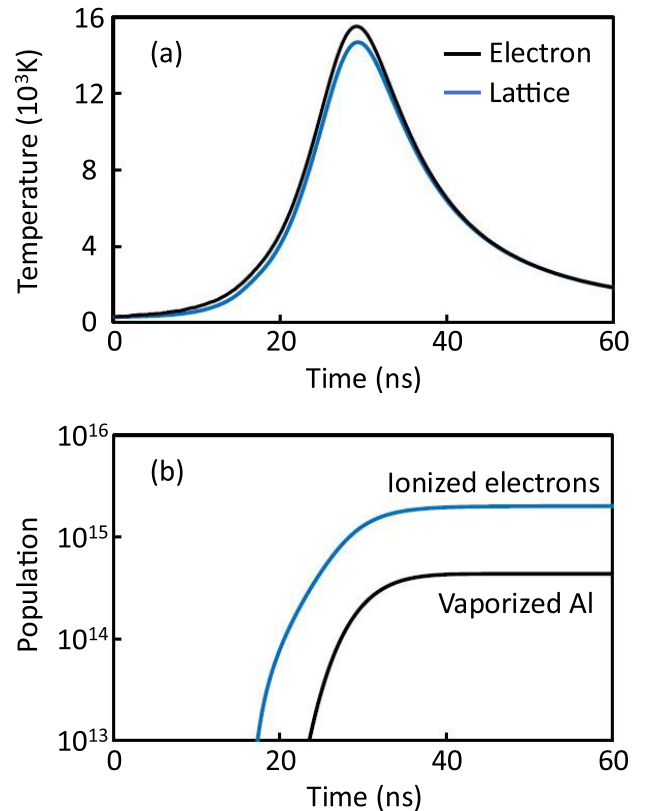


Fig. 10. (a) Conduction electron and lattice temperatures T_e and T_l versus time. (b) Ionized electron population and vaporized Aluminum atoms versus time.

Simulations show that the electron plasma inside the Aluminum increases in total energy and effective temperature due to free-carrier absorption and electron–electron scattering. Past the laser pulse peak, electron–phonon scattering becomes the dominant effect, resulting in the cooling of the electron plasma. The starting distribution (Fermi–Dirac distribution at 300 K with chemical potential at 10 eV above the conduction band edge) begins to change at 10 ns, even prior to any

TABLE IV
INPUT PARAMETERS FOR LASER-AL(SOLID) INTERACTION

<u>Laser</u> Photon energy Pulse duration Beam diameter	$\hbar\nu = 1.17\text{meV}$ $\tau = 10\text{ns}$ $d = 220\mu\text{m}$
<u>Aluminium</u> Specific heat Density Boiling temperature Enthalpy of vaporation Fröhlich electron-LO phonon coupling LO phonon frequency Electron effective mass Optical dipole matrix element Carrier-carrier scattering rate Carrier-phonon scattering rate Heat conduction rate Work function	$\sigma_{Al} = 9 \times 10^2 \text{J/kg/K}$ $\rho_{AL} = 2.7 \times 10^3 \text{kg/m}^3$ $T_b = 2743\text{K}$ $L_{Al} = 284\text{kJ/mole}$ $\hbar G_q = 6.3\text{meV}$ $\hbar\omega_{LO} = 36.4\text{meV}$ $m_c = 0.97m_0$ $\wp = e \times 0.2\text{nm}$ $\gamma_e = 5 \times 10^{13} \text{s}^{-1}$ $\gamma_l = 10^{12} \text{s}^{-1}$ $\gamma_c = 10^{11} \text{s}^{-1}$ $W = 4.08\text{eV}$

TABLE V
INPUT PARAMETERS FOR RATE EQUATION

Electron impact excitation $R_{Ar(4s)} : \beta_{Ar(4s)} = 3.15 \times 10^{-22} \text{m}^2 \text{eV}^{-1}, \quad \varepsilon_{Ar(4s)} = 11.7\text{eV}$ $R_{Ar(4p)} : \beta_{Ar(4p)} = 3.15 \times 10^{-22} \text{m}^2 \text{eV}^{-1}, \quad \varepsilon_{Ar(4s)} = 13.2\text{eV}$
Photoionization $R_{Ar(*+)} = 3 \times 10^{-26} \text{m}^2, \quad \varepsilon_{Al(+)} = 6.0\text{eV}$
Recombination & spontaneous emission $\gamma_{4s}^{sp} = \gamma_{4p}^{sp} = 10^6 \text{s}^{-1}$ $R_{Al(+)} = 5.6 \times 10^4 \text{m}^{-3} \text{s}^{-1}$

appreciable laser intensity. From 20 ns onward, the laser intensity increases rapidly and free-carrier absorption begins to dominate. Plasma heating continues until around 30 ns at the peak of the laser pulse, after which there is a cooling of the electron distribution with energy transfer from the plasma to the lattice by electron–phonon scattering. Somewhere around the peak of the laser pulse is a small decrease in total electron population because of thermionic emission. The high

temperatures lead to thermionic emission of electrons and vaporization of Aluminum.

To connect to experimental data, the next step is to model the excitations and optical emission from Argon atoms (buffer gas) using a highly condensed sketch of the relevant energy levels and optical transitions (Fig. 11). These are the same processes that would be included in a PIC simulation if we were modeling a laser-triggered switch.

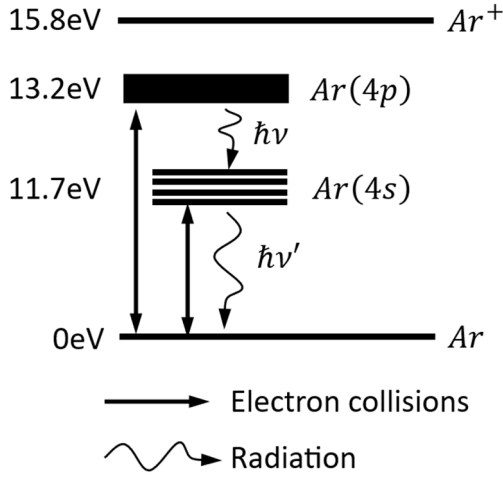
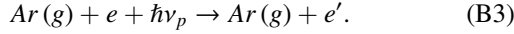
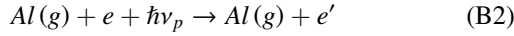
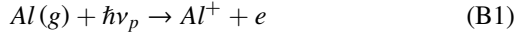
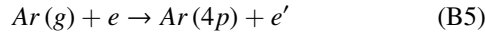
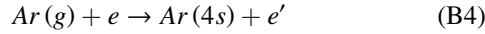


Fig. 11. Electronic levels and relevant optical transitions in the Argon atom.

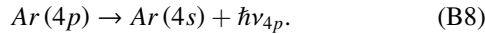
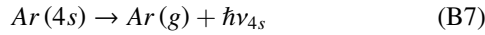
Laser photons not lost to free-carrier absorption give rise to photoionization and inverse *Bremsstrahlung*



Electrons from thermionic emission and photoionization can excite Argon atoms or recombine with Aluminum ions as follows:



Lastly, the excited Argon atoms will relax via spontaneous emission



We take as input the densities and energies of thermionically emitted electrons and vaporized Aluminum atoms shown in Fig. 10. The following rate equations are solved for the ionization and recombination processes:

$$\frac{dN_e}{dt} = R_{Al(+)}N_{Al(g)}S_p - R_{Al(g)}N_{Al(+)}N_e \quad (B9)$$

$$\frac{dN_{Al(g)}}{dt} = -R_{Al(+)}N_{Al(g)}S_p + R_{Al(g)}N_{Al(+)}N_e \quad (B10)$$

$$\frac{dN_{Al(+)}}{dt} = R_{Al(+)}N_{Al(g)}S_p - R_{Al(g)}N_{Al(+)}N_e. \quad (B11)$$

For the Argon atoms, there are electron impact excitation and radiative recombination processes described by the rate equations

$$\begin{aligned} \frac{dN_{Ar(g)}}{dt} = & -R_{Ar(4s)}N_{Ar(g)}N_e - R_{Ar(4p)}N_{Ar(g)}N_e \\ & + \gamma_{4s}^{sp}N_{Ar(4s)} + \gamma_{4p}^{sp}N_{Ar(4p)} \end{aligned} \quad (B12)$$

$$\frac{dN_{Ar(4s)}}{dt} = R_{Ar(4s)}N_{Al(g)}N_e - \gamma_{4s}^{sp}N_{Ar(4s)} \quad (B13)$$

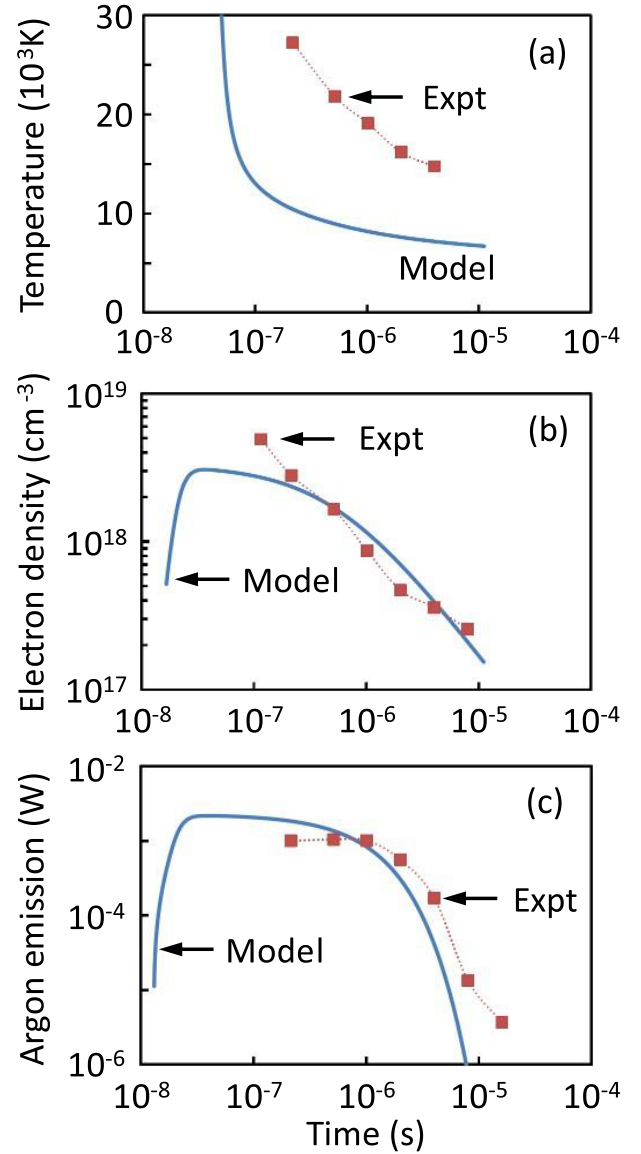


Fig. 12. Outside the Aluminum target: (a) electron temperature, (b) electron density, and (c) Argon emission versus time. The dots from the experiment and the curves are calculated results. For (c), the y-axis scale is for the 4p-4s transition.

$$\frac{dN_{Ar(4p)}}{dt} = R_{Ar(4p)}N_{Al(g)}N_e - \gamma_{4p}^{sp}N_{Ar(4p)}. \quad (B14)$$

Lastly, the electron plasma energy outside the Aluminum target is determined by imposing energy conservation

$$\begin{aligned} \frac{dN_e kT}{dt} = & \eta S_p \hbar\nu_p - \varepsilon_{Al(+)}R_{Al(+)}N_{Ar(g)}S_p - \frac{3}{2}k_B T R_{Al(g)}N_{Al(+)}N_e \\ & - (\varepsilon_{Ar(4s)}R_{Ar(4s)} - \varepsilon_{Ar(4p)}R_{Ar(4p)})N_{Ar(g)}N_e \end{aligned} \quad (B15)$$

where the first term on the right-hand side is the laser energy transferred to the electron plasma, the 2nd and 3rd accounts for ionization and recombination in gaseous Aluminum, the 2nd line accounts for energy lost to electron impact excitation

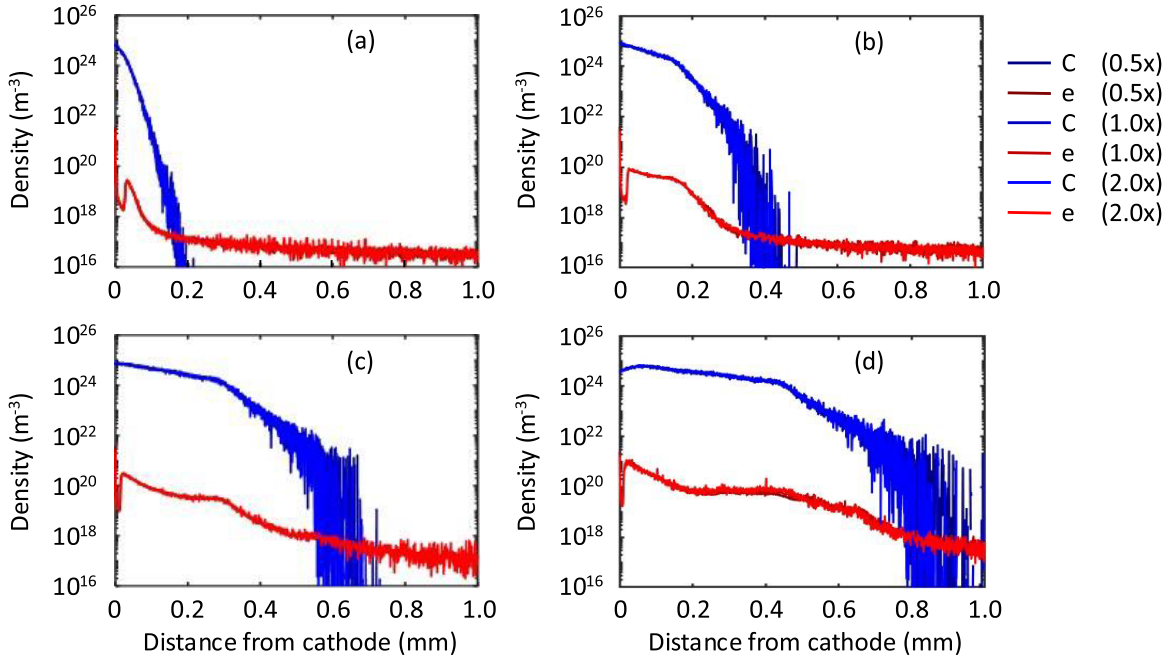


Fig. 13. Carbon atom (blues) and electron (reds) densities across the entire 1 mm gap. The times are (a) 25, (b) 50, (c) 75, and (d) 100 ns, with results averaged over a 0.01 ns window. For each time, the three different curves are shown that use spatial discretization and temporal step sizes of $dx = 0.1175 \mu\text{m}$, $dt = 0.5 \cdot 10^{-13} \text{ s}$, $N_p = 200$ (dark blue and dark red, labeled as 0.5x), $dx = 0.2350 \mu\text{m}$, $dt = 1.0 \cdot 10^{-13} \text{ s}$, $N_p = 100$ (blue and red, labeled as 1x), $dx = 0.4700 \mu\text{m}$, $dt = 2.0 \cdot 10^{-13} \text{ s}$, and $N_p = 50$ (light blue and light red, labeled as 2x). Note that the curves for each numerical discretization almost entirely overlap.

of Argon. The cross sections and rates have the general form

$$R_X = \sqrt{\frac{8}{\pi m_x}} \beta_X \sqrt{kT} (\epsilon_X + 2kT) \exp\left(-\frac{\epsilon_X}{kT}\right) \quad (\text{B16})$$

where Table V gives the parameter values

Fig. 12 shows the results of numerically solving the rate equations compared to experimental data (courtesy of Xianglei Mao at LBNL).

There is encouraging closeness to experimental data, sufficient for the PIC simulations. Further improvements would include the incorporation of the effects from the ionization of Ar, excitation of Al, and the resulting emission spectrum, and expansion of the plasma. An important input parameter is the fraction of the laser energy that went into the ionization of Al. This free parameter could be eliminated through a combination of the light-matter interaction and the free-space interactions into a single model.

APPENDIX C

This Appendix provides convergence results for the PIC-DSMC simulations used, demonstrating that our results are consistent across different mesh sizes, time step sizes, and computational particle counts per element. In particular, Fig. 13 shows carbon and electron densities across the vacuum gap for four different times and three different discretization choices for each time, revealing quantitatively similar results regardless of the numerical parameters used. To augment the data provided in Fig. 5, these results are shown at several early times to show the evolution of these densities as the plasma expands across the gap and closes the switch.

DATA AVAILABILITY

The data that support the findings of this study are available from the corresponding author upon reasonable request.

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