PAPER

Laser-plasma induced shock waves in micro shock tubes

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Abstract

Shock wave generation with help of lasers or shock tubes, for example, is a common subject at macroscopic scale. On the other hand, the current tendency towards smaller scales becomes more and more important. In particular, shock waves in small channels or tubes with sub-mm or micron-sized diameter have attracted much interest within the last years. But downscaling of shock wave effects such as pressure decrease and Mach number attenuation during propagation, viscous and heat effects, laminar and turbulent flow is not necessarily straightforward from macro to micro or even nano range. Although several theoretical investigations are available in this new field, there is a strong demand on experiments, which are mostly missing due to the lack of suitable methods. The present work introduces a novel method for the generation of shock waves at microscale, namely laser-induced micro shock waves (LIMS). The LIMS method applies a femtosecond laser to induce an optical breakdown in a thin aluminum target located at the entrance of a micro tube. Subsequently, a shock wave is launched by the high pressure aluminum plasma and starts propagating into the tube. The topical work presents, for the first time, experimental investigations on direct micro shock wave generation and propagation at well-defined conditions in micron-sized tubes. They are performed for different conditions and tubes down to 50 μm diameter. Different from previous shock wave investigations in the mm-diameter range that involve pressure transducers, the present work applies non-contact measurements by optical methods. The experiments are supported by additional simulations. A one-dimensional numerical hydrocode is applied to simulate the shock wave generation process. Further propagation of the shock in a micro tube is analyzed by solving two-dimensional Navier–Stokes equations. Both simulations agree well with the experimental results.

Nomenclature

| Symbol | Description |
|--------|-------------|
| τₐ     | laser pulse duration (fs-laser) |
| Eₐ     | laser pulse energy (fs-laser)   |
| I      | laser intensity (fs-laser)      |
| φ      | laser focal spot diameter (fs-laser) |
| f-num  | f-number                              |
| τ₏     | shock buildup time               |
| γ      | heat capacity ratio              |
| Re     | Reynolds number                   |
| P, p   | pressure                           |
| μ      | dynamic viscosity                 |
| ν      | kinematic viscosity               |

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Investigation of shock waves has a rather long history. This comprises shock waves generated by (large) shock tubes (with up to several m diameters) and by intense laser pulses. Shock waves were also driven by detonations, pinch plasmas (typical size is a couple of mm or cm) and other devices, or on a much smaller scale e.g. by the piezoelectric effect (typical size is mm). However, we intend to discuss here a different scheme of shock wave generation and propagation: the laser-plasma induced micro shocks (LIMS), which propagate in a confined environment at micro scale (sub-mm towards dozens of μm). Therefore, the current work includes two major topics in the shock wave research: laser induced shocks and shock miniaturization. The combination of the two can be considered as new in this research domain.

In modern era of high power laser, a new branch of physics in addition to the classical shock wave theories starts to play a role (e.g. [1–3]). Here, a single pulse or multiple intense laser pulses are usually focused to a small volume onto a solid target, and then a shock wave occurs and propagates into the target. Related experiments are performed to investigate high energy density physics (HEDP) [4] with special attention to the equation of state (EOS) [5], inertial confinement fusion (ICF) [6] and astrophysics [7]. Experiments in HEDP rely on very high concentration of energy in a small volume, which usually leads to high density, temperature and pressure. HEDP may already happen in the laser pulse duration, especially for the case of ultrashort laser pulses. Consecutive ablation of the irradiated surface also has great importance (particularly for ICF) in building high pressure and density. As an example, in the first experiments of Sigl [8] in 1974, pressure in excess of one Mbar was achieved. Since the last two decades, pressure in the range of multiple GBar became quite common [9]. Due to the short duration of the laser pulse, this pressure is produced rather fast [10]. One step further, if the laser intensity is big enough, the light pressure itself (ponderomotive pressure) may contribute to the shock onset along with the ion and electron pressure. For the laser intensity of $10^{16}$ W cm$^{-2}$, the ponderomotive pressure is in the Mbar range. For the intensity of $10^{19}$ W cm$^{-2}$, the corresponding pressure will reach the Gbar range. This phenomenon was discussed by Lebedev as early as in 1901 [11]!

When the high power laser systems based on ultrashort pulses became available, they have been applied to investigate shock waves (e.g. [3, 12]). Nowadays, femtosecond (fs) and picosecond (ps) laser-induced shocks already find practical applications, especially in micro technology e.g. the removal of thin layers [12]. Jiang and Takayama [13] have shown that the laser-plasma induced shock wave is approximately a blast wave created by a point-like explosion, which can be described by self-similarity methods [14]. More recently, fundamental investigations on unconfined propagation of shock waves generated with fs-lasers have also been extended to microscale [15, 16].

The aforementioned examples and the advancing micro and nanotechnology do strongly indicate the importance to extend shock wave investigations to the micron range as a multidisciplinary topic. As this work concentrates mostly on the shock waves which propagate in a ‘confined’ geometry, such as provided by tube, we may define different ranges according to the tube diameter (or the corresponding hydraulic diameter, see below). Within this work shock waves generated in tubes with diameters larger than one mm are attributed to
‘macro range’, whereas if it is sub-mm down to a couple of microns, these are referred to ‘micro range’. Of course, the flow properties of shock waves on micro scale may differ from those on macroscale. This is due to the increased significance of scaling effects which are functions of either 1/D or 1/D^2 (where D is a characteristic length, such as the hydraulic diameter) [17]. Hence, friction becomes non-negligible, heat conduction leads to non-adiabatic flows, and small inhomogeneities may play a more significant role. Such conditions that are mostly idealized within investigations on macroscopic fluid dynamics, now cannot be neglected anymore.

Recent years, Brouillette [17] has introduced a dimensionless scaling factor Sc = Re · D/(4L) (L denotes the distance between the shock and the contact surface) to characterize the scaling of shock flow. Sc indicates that lower initial pressure and smaller hydraulic diameter shall have the same scaling effects. Zeitoun [18] has simulated the flows in micro shock tubes by computing two-dimensional Navier–Stokes equations combined with slip conditions. More recent works on micro shock waves are presented in e.g. [19, 20].

However, experimental works at even smaller scales are still largely missing due to the limitations in micro-processing and measurement technologies. Actually, to the best of our knowledge, there is only one experiment with D < 100 μm [21, 22]. So far, the shock wave was initially generated in a larger conventional tube and then transferred into an inserted narrower tube.

Generally speaking, shock waves are of great interest in the fundamental research as well as applications, e.g. explosions and high-temperature physics [23], medical technology [24], rare metal recycling [25], aerospace engineering [26–28], electrohydrodynamics [29] and cold gas coating [30]. Beside the interest in fundamental researches, the current work may be of relevance for microfluidics [31] or even nanofluidics [32]. Of course, similar to other examples in the history of science, currently unforeseen applications may emerge later on. The current work also presents, for the first time, experimental investigations on direct micro shock wave generation and propagation in micro tubes at well-defined conditions. Thus it can avoid difficulties that may occur at an abrupt area change from a larger tube to a smaller one, as mentioned above.

Micro and macro shock flows can have even more fundamental differences, due to different governing regimes of statistical mechanics and continuum mechanics. The validity of continuum mechanics can be indicated by the well-known Knudsen number Kn = Λ/D with Λ as collisional mean free path. The Knudsen number for our experimental conditions is in the range 0.001 < Kn < 0.01 for the smallest tube with D = 50 μm, because Λ ≈ 68 nm for atmospheric conditions [33]. The Kn here is also comparable with the one in the theoretical work [34]. Therefore, as we will see, although LIMS in the current work is different than macro shock waves due to dissipative effects, continuum mechanics still applies here. If the hydraulic diameter of the shock flow is further downscaled from μm to sub-μm scale, Kn could be close to 1, thus molecular physics and statistical mechanics should be considered. Consequently, there is a strong demand on new concepts and developments in the miniaturization of shock tubes.

The objective of the present work is to combine the major issues addressed above: downsizing of shock waves in a miniaturized tube (within the present work the term ‘tube’ is used for any kind of tube, capillary, channel etc, independently on its size) and the development of a novel method for shock generation by applying laser plasma physics. The main goal of this work is to open up a new scheme to investigate the fundamentals of micro shock waves, thus it shall not be regarded as a complete investigation of all its details.

This article is ordered as follows. Section 2 describes the LIMS method. In section 3, laser plasma generation and the onset of the shock wave using air or water as the medium for shock propagation is modeled by hydrodynamic simulations. Section 4 concentrates on the experimental setup. In section 5, the experimental results of shock propagation are described, analyzed and compared to two-dimensional Navier–Stokes computation. Differences to the classical shock tubes, scaling effects, core flow expansion behind the shock wave as well as the development of the wall boundary layer will be discussed.

2. The LIMS Method

In contrast to shock wave generation with macroscopic tubes, which have a high pressure part and a low pressure section, the present LIMS method makes use of a tube filled with an initially homogenous fluid in its whole volume. The principle of the LIMS is shown in figure 1. The set-up consists of a 1 mm thick glass plate located very closely to the entrance of the tube, namely at a distance which is at least one order of magnitude smaller than the tube diameter, even for the smallest tube. Thus, the laser-initiated shock is almost completely lanced into the tube. The present work includes experiments, where the opposite side of the tube was open.

An intense laser pulse is focused through the glass plate onto a thin layer of aluminum deposited on the rear side of the plate (i.e., the target; note: such thin layers can be partly oxidized), where it generates almost instantaneously a laser-produced plasma (LPP). The sudden occurrence of a high-pressure high-temperature LPP represents an extreme non-equilibrium, which runs down in the emission of a shock wave. This LPP acts as the driver for a shock wave, which propagates into a tube positioned in the immediate vicinity (less than 5 μm
from the target). Within the present work, much care is taken to provide clear and well-defined conditions for the shock wave generation, but the shock wave strength itself is of minor importance meanwhile.

To achieve this goal, the LPP should be set on nearly instantaneously as a ‘δ-pulse’, which could be described approximately by the mathematical δ-distribution. This is fulfilled, if the corresponding time scale of the driver process is much shorter than that of the shock wave buildup. In such case, details of its generation and evolution are not important for the shock wave propagation. Moreover, the LPP should act as a homogenous and planar driver. These conditions are fulfilled, if the (initial) thickness $d_{\text{LPP}}$ (during the onset of the shock) is much smaller than the lateral width of the driver. To achieve the desired conditions, the laser pulse duration $\tau_L$ has to be rather short, namely in the sub-ps or fs range. Then, according to hydrocode simulations (see section 3), $\tau_L$ is much shorter than the time $\tau_i$ necessary to set up the shock wave.

Here it should be remarked that shock waves may be also directly generated in the medium, where its propagation will be investigated (e.g., in air or water). However, this method may inclose problems in providing ill-defined conditions for shock wave onset and its propagation. As an example, LPP generation in a water-filled small capillary leads to cavitation bubbles with a rather complicate evolution and propagation with many reflections at the capillary walls [35, 36].

For the present work, linearly polarized pulses of a Titanium:Sapphire laser system with a wavelength of $\lambda_L = 775$ nm, adjustable pulse energy $E_L$ (maximum $E_L = 1$ mJ, but typically $E_L = 20$ $\mu$J in the experiments) and a pulse width $\tau_L = 150$ fs (FWHM) were focused with an achronomic lens of rather long focal length $f$ at normal incidence. This yields a large focus but still provides a sufficiently high intensity $I > 10^{13}$ $\text{W cm}^{-2}$ to induce an appropriate LPP for subsequent shock wave generation in front of the tube. Preliminary experiments were performed with different lenses and thus different values of $f$ and with intensities $I$ well beyond the plasma formation threshold of Al ($I_{\text{Al}} = 2 \times 10^{12}$ $\text{W cm}^{-2}$). Thereby, the laser intensity was kept well below the ionization threshold of air ($I_{\text{air}} = 5 \times 10^{14}$ $\text{W cm}^{-2}$) and also close to the threshold of optical breakdown (or damage) in glass ($I_{\text{glass}} = 2 \times 10^{13}$ $\text{W cm}^{-2}$). All values given in brackets are deduced experimentally for conditions of the present work. To generate shock waves sufficiently strong for the propagation measurements (see section 5), $I$ is chosen inevitably close to $I_{\text{glass}}$ (see table 1). However, the optimization of shock wave strength was not in the focus of present work. It should be mentioned that even $I > I_{\text{glass}}$ may be applicable as long as $I/I_{\text{glass}}$ is not too large. This is because all intensities within the present work represent the peak values in spatial

Table 1. Parameters for shock wave generation. ‘Planar’ and ‘point’ correspond to the geometry of the LPP as shock driver. Although laser pulse energy $E$ differs for the different geometries, laser intensity $I$ is kept constant and thus the initial pressure $P$ of the LPP remains constant (with exception of the point-like geometry: here $E$ is the same as in planar geometry and thus $I$ and $P$ are an order of magnitude larger).

| $D$ ($\mu$m) | No tube | Planar | Point |
|-------------|---------|--------|-------|
| $f$ (mm)    | 150     | 150    | 300   | 750   | 40 |
| $f$-num     | 30      | 30     | 61    | 153   | 40 |
| $\phi$ (ºm) | 21      | 21     | 42    | 106   | 6  |
| $E_L$ ($\mu$J) | 20     | 20     | 120   | 740   | 20 |
| $I$ ($10^{13}$ $\text{W cm}^{-2}$) | 2       | 2      | 2     | 2     | 30 |
and temporal profiles of the laser beam (for a Gaussian profile this intensity can be simply estimated as 
\[ I = 0.65 \cdot \frac{E}{(\tau/\phi^2 \pi/4)}, \]
where \( \phi \) is the FWHM-diameter of the focal spot). Therefore, even for \( I > I_{\text{glass}} \), the ionization threshold of glass would be exceeded only in the relatively tiny central part of the beam during the pulse maximum, and this does not affect the LPP within the thin Al-layer too much. In shots without the Al layer, shock wave generation was absent even for the highest supplied intensity (see table 1). Consequently, (significant) plasma formation is restricted to the thin Al-layer, which later on drives the shock wave into the tube. After each shot, the target is shifted, so that the next laser pulse is focused to a fresh spot on the target (figure 2). Different layers (at least partly oxidized) of different thickness between \( d_{\text{Al}} = 30 \) and 100 nm have been investigated to determine a suitable value of \( d_{\text{Al}} \) with respect to LPP driver action.

Proceeding to the experiments, an additional reflectivity and transmission measurement has been carried out illuminating a target with layer thickness \( d_{\text{Al}} = 50 \) nm (see figure 2). A reflectivity \( R = 65\% \) (measured at \( I = 2 \times 10^{13} \text{ W cm}^{-2} \)) was detected with help a thermopile detector (Coherent, LabMax-TOP). The detector was also applied to measure the transmission \( T \) at the same conditions, but the transmitted energy appeared to be smaller than the sensitivity limit of the detector. Applying a photodiode, the transmission could be determined as \( T = 4.8\% \) (measured at somewhat lower \( I \)). Thus, the penetration depth is \( \lambda \approx 16.5 \) nm (of course, this value changes significantly with time during the laser plasma interaction). Due to oxidation, this initial value is larger than the value of pure Al (\( \lambda \approx 7.1-7.5 \) nm for Al [37]), but it is still small enough to provide strong absorption of the energy supplied to the thin Al-layer (absorption \( \lambda = 30\% \); hence 6 \( \mu \)l out of the 20 \( \mu \)l incident onto the target is deposited and used for shock wave generation; A similar value is found in the hydrocode simulations described in section 3).

Assisted by the very fast nonlinear heat wave within the thin layer, the LPP is almost homogeneously heated and ionized. Thus the electron pressure \( p_e \) is nearly constant all over the LPP. These initial plasma conditions may be quickly estimated from a ‘\( \delta \)-pulse model’ based on the work of Zeldovich and Raizer [23], which describes the initial plasma conditions for ultrashort intense laser pulses [38, 39]. For the typical experimental value of \( I = 2 \times 10^{13} \text{ W cm}^{-2} \) in the present experiments (peak value, see table 1), the initial electron temperature \( T_e \) reaches several eV, the average ionization degree of the plasma is approximately 3 and thus the electron number density amounts to \( n_e \approx 2 \times 10^{23} \text{ cm}^{-3} \), which constitutes hundred times of the critical density \( n_c \). This yields an initial electron pressure \( p_e \) of several MBars. During the laser pulse itself, due to the large inertia of the ions, the electrons and ions are not in equilibrium, since the ion temperature \( T_i \) is still close to room temperature. Only after a few ps both temperatures may equalize to a common temperature \( T \), if the characteristic electron–ion relaxation time is sufficiently small, which is also a function of the electron number density.

Within a few ps, the Al-LPP begins to expand significantly and pushes the ambient air (or water) inside the tube. A first estimate of the relevant acoustic impedance at the solid density target/fluid interface may be evaluated from basic principles as described in textbooks (e.g., see [40]). Applying the sound velocity \( u_{\text{ct}} \), the pressure \( P \), and the mass density \( \rho \) of the LPP and air, the acoustic intensities \( I_{p}, I_{A} \) and impedance values \( Z_p, Z_A \) of the LPP and air, respectively, can be estimated. For the LPP, a \( \delta \)-pulse model evaluation yields \( u_{\text{ct}} \approx 10^4 \text{ m s}^{-1} \) and \( P \approx 10^7 \text{ Mbar} \) (prior to expansion, \( \rho_b \) is assumed to be the same as that of solid Al; for air, standard conditions may be assumed; here and in the following, the indices ‘\( p \)’ and ‘A’ indicate plasma and air).
As a result, the wave intensity transmitted into air is given by

$$I_A = 4I_P \frac{Z_p Z_A}{(Z_p + Z_A)^2}$$

(1)

and the transmitted pressure into air yields

$$P_A = I_A / \mu_A \approx 1.8 \times 10^4 \text{ bar.}$$

(2)

For an intensity ($I_P \approx 10^{12} \text{ W cm}^{-2}$ in the aluminum layer, we get impedance values $Z_p = \rho_p \cdot \nu_{ap} = 2.7 \times 10^7 \text{ kg m}^{-2} \text{s}^{-1}$, $Z_A = 428 \text{ kg m}^{-2} \text{s}^{-1}$).

Thus, as expected, the initial driving pressure of the LIMS (i.e., $P_A$) is significantly smaller than that of the LPP. Nevertheless, this very rough estimate promises strongly pushed shocks. Due to rapid evolution of the plasma, the impedance mismatch is eased later on. Therefore, no extra effort is taken yet to improve the acoustical impedance within the present work e.g. by using special targets, which may be a subject of future work. Note that the expected impedance adaptation is confirmed by the hydro simulations of plasma dynamics in the aluminum–air compound described in section 3.

Figure 2 shows the image of a target after exposure with one fs-laser pulse. The diameter of the ablated region fits well to the diameter of the tube cross section and thus supports the assumption of a planar shock wave. It leaves an undisturbed glass surface behind. A similar statement can be made for bigger tubes analyzing a similar microscope image of the target. The smooth surface outside of the ablation region shows the condition prior to the exposure.

3. Hydrocode simulations of shock onset

To get a principal insight in the short-pulse laser-target heating and the subsequent propagation of a shock wave in the adjacent medium, one-dimensional simulations with the MULTI-fs code [41] were performed. This Lagrangian hydrodynamic code with multi-group radiation transport simulates the laser pulse propagation in the plasma region up to the critical surface by solving the wave equation. As the result, light reflection in plane geometry is modeled correctly and provides realistic absorption values. Hence, we avoid an overestimation of the gasdynamic pressure by excessive absorption in the aluminum layer.

Since ion and electron fluids in short-pulse-driven plasmas may be far from thermodynamic equilibrium, the code implies separate EOS tables for both species. The implementation of this feature and test simulations with aluminum targets were reported in [42]. In the present analysis, EOS data for aluminum were calculated with help of FEOS [43, 44], a code based on the QEOS model [45] and upgraded by Frankfurt University. It includes the so-called soft-sphere approximation [46], which avoids overestimated plasma pressures in the two-phase region up to the critical point. Simulations were performed with two different propagation media: dry air and water.

EOS tables are taken from SESAME library [47, 48]. Since this library offers only the general table 5030.301 for dry air, which contains cumulative pressures and specific internal energies, the electron EOS was evaluated by subtraction of an ideal gas EOS for the ion fluid from this table. In general, such an assumption is justified for sufficiently high temperatures. This proves to be especially significant for the electron tables. The situation becomes relaxed in an air medium because of its low density. At temperatures above several electron volts, which are typical for the shock wave region, the deviations in the approximate EOS tables should be tolerable. In case of a water medium, SESAME provides separate EOS tables 7154.304 and 7154.305 (ion EOS containing zero point data but not the cold curve). Comparing these tables with approximate ones generated as described above for air, we may confirm our previous argumentation concerning the influence of plasma temperature. Note also, that the ideal gas pressure isotherms at normal density are now surprisingly close to the corresponding curves from ion EOS 7154.305 even at temperatures far below an electronvolt.

Since the MULTI-fs code solves equations for electron and ion internal energies, inverse EOS tables are used, which contain pressure and temperature data as functions of plasma mass density and specific internal energy. The knowledge of temperatures is required for the evaluation of transport processes in the plasma. Electron collision frequency in the classical plasma parameter range as well as in the range of dominating electron–phonon interaction depend on temperature and determine absorption and electron–ion energy relaxation. Finally, electron thermal conduction is taken into account by means of a harmonic law. It stands for the limitation of an unphysically high collisional Spitzer conductivity in steep temperature gradients by the natural thermal electron flux. Since our study has shown gentle gradients, a flux limiter of 0.6 was used, which favors the classical conduction process. Opacity coefficients for radiation transport simulations were calculated with help of SNOP [49], which is a stationary non-equilibrium opacity code. Since x-ray emission and re-absorption appeared to be weak in the performed simulations because of low temperatures therein, opacities for only the...
nitrogen component were used in air. In water, we applied opacities for oxygen. Charge numbers in aluminum were gained from FEOS simulations, the corresponding values in air and water were evaluated with SNOP.

In accordance with experimental conditions, a laser pulse at wavelength 0.775 μm, FWHM duration of 150 fs with a sin-squared intensity envelope and a peak intensity of 2 × 10^{13} W cm^{-2} were assumed. The transparent glass support for the 50 nm aluminum layer was mimicked by the boundary condition of zero flow velocity on the laser-illuminated aluminum boundary. Typically, an amount of 26% of the laser pulse energy is absorbed by the aluminum layer. This result is close to the experimentally observed value. Simulations also show that the electron number density is kept overcritical, which confirms the measured very low transparency (see section 2).

Spatial profiles of mass density, ion and electron pressures, ion and electron temperatures and flow velocity as functions of the shock propagation distance x in case of an air medium.

Figure 3. Mass density ρ, ion and electron pressures $p_i, p_e$, ion and electron temperatures $T_i, T_e$ and flow velocity $v$ as functions of the shock propagation distance x in case of an air medium.

In addition, we realize a slight deceleration of the shock front in air already at this early propagation stage (see, e.g., panel (b)). Since the hydrodynamic equations of the code contain only an artificial viscosity term to broaden the shock front over a couple of mass cells, there is no feasible mechanism for this effect except some downstream expansion flow out of this shock. A relaxation wave appears also in the aluminum plasma behind the pressure maximum at the interface with air.

The dynamics of the shock wave induced by the gasdynamic pressure in the laser-heated aluminum layer becomes somewhat more vivid in figures 4–6, which demonstrate plasma profiles as functions of mass cell numbers, since one can easily distinguish the boundary between aluminum plasma and air in all plots. The left 100 cells on the simulation grid are filled with aluminum.

Analyzing figure 4, we focus the discussion on shock generation in the aluminum layer. Shock breakout into air appears at approximately 1 ps (red curves). At this time, electron pressure reaches approximately 1 Mbar in the aluminum plasma, while an electron temperature of about 5 eV will be established there (see panels (e), (f)). The subsequent expansion of the aluminum plasma into the air volume drives a strong relaxation wave as mentioned above. This becomes evident from, e.g., figure 4(a).
Figure 5 demonstrates the proceeding shock wave evolution at later times, $t \gg 10$ ps, like in figure 3. A stable shock wave in air is formed, with a slowly decreasing ion pressure amplitude in the range of a couple of kbars. Obviously, the initially large pressure discrepancy in aluminum and air, which represents the strong impedance mismatch, balances after some time because of the aluminum layer expansion. The pressure wave heats the air medium to about 10 eV. Such temperatures imply values for the sound velocity between 6 and 9 km s$^{-1}$. Consequently, Mach numbers between 2 and 3 can be extracted from the simulated data, which agrees well with the experimental results (see section 5). Looking at the electron pressure and temperature plots, one can observe an area adjacent to the aluminum layer, which is probably affected by heat conduction. This effect is also visible in figure 4(f). Radiative transport cannot be responsible for it, since a simulation without such transport did not change the shown results. Electron temperatures of about an electronvolt are found here. In the remaining part of compressed air, the electron fluid is heated only marginally, unlike the ion component.

From the analysis of the present numerical results, one may suggest that the shock will become considerably weaker at times $t \gg 1$ ns (as it is shown in section 5.3).
Since temperature also decreases thereby, the corresponding Mach numbers may remain in the earlier estimated range.

Finally, we discuss the evolution of the electron number density in the air volume, depicted in figure 6. It appears to be about two orders of magnitude below the numbers in aluminum, which is mainly a consequence of the much lower density in air (see panel (a)). Therefore, the ionization degree becomes very low at later times because of reduced collisional ionization. This fact is demonstrated in figure 6(b).

For comparison, simulations with a water medium behind the aluminum layer were performed. Applying the same laser parameters as before, an absorption coefficient of 0.26 is obtained again. The resulting spatial profiles of plasma parameters, however, significantly differ from those observed with an air medium. First of all, the shock propagation proceeds much slower in water, as shown in figure 7. Flow velocities are an order of magnitude smaller than in air because of the much higher mass density in the water medium. Since, in addition, sound velocity in water is several times larger than in air ($u_a = 1.54 \text{ km/s}$ at normal conditions), smaller Mach numbers are accessible.

Concerning ion pressure characteristics (see figures 3 and 7(c)), we observe significantly higher values in water—about 20 times more at $t = 1 \text{ ns}$ than in air, which is a consequence of the much better impedance matching.

Comparing the temperatures of electrons and ions, similar profiles can be observed in the aluminum plasma as well as in the water volume, in contrast to the case with the air medium. Now, the ions are heated only to a sub-electronvolt level, and both fluids are thermodynamically in quasi-equilibrium except a small deviation in the region of maximum ion pressure (see figures 7(d) and (f)).
Since the relatively high mass density of water hampers a remarkable expansion of the aluminum layer, the interface between driver and propagation medium cannot be distinguished so easily as in the case with air. Therefore, we turn—as in case of an air medium—to the presentation of the simulated plasma parameters as functions of the mass cell number (see figures 8–10), which directly exhibits the aluminum–water boundary.

Figure 8 demonstrates plasma evolution in aluminum and the shock propagation in water during the first 10 ps after the onset of laser illumination. As expected, the shock breakout from the aluminum layer appears again at approximately 1 ps. Reduced flow velocities in the aluminum layer in comparison to the case with an adjacent air volume follow from figure 8(b). Electrons and ions reach thermodynamic equilibrium at times $t > 10$ ps (see panels (d) and (f)). Obviously, the reduced aluminum expansion helps to keep an efficient electron–ion relaxation process.

Plasma profiles as functions of the mass cell number at times $t \gg 10$ ps are depicted in figure 9. An interesting feature is shown in panel (a), where the mass density in compressed water becomes larger than in the
aluminum layer for curves with \( t \geq 50 \) ps. In addition, a strong drop of density is observed in the water downstream region, which may be caused by an intense expansion flow, as already discussed for the shock propagation in air. Because of the higher ion pressures of shocks in water, in comparison with the air case, the density decline becomes more pronounced. The remarkable decrease of the resulting flow velocity in this region (see panel (b)) supports such an interpretation. Also, the stronger damping of the shock (see panel (c)), as compared to the shock decrease in air, may be understood as a further indication.

Analyzing pressures and temperatures in the electron fluid (see figures 9(e) and (f), respectively), visible profiles in the whole volume of compressed water can be obtained now, in contrast to the situation in air. The reason for that is demonstrated in figure 10. Electron number density becomes maximum in the shock region and gains high values above the numbers of the corresponding Fermi gas (panel (a), black curve). Furthermore, the large number of free electrons is responsible for an ionization degree between one and two despite the relatively low temperature. If we would additionally account for hydrogen atoms in the heated water, these values should be even higher.

In summary, we would like to remind that the limited quality of the applied EOS especially in the low-temperature range may distort the quantitative output of the simulations at some extent. This is different to standard laser-plasma simulations with higher intensities or with long-term illumination. Nevertheless, the plasma parameters in the shock wave region at a final instant around a nanosecond in our MULTI-fs simulations may be regarded as a planar input for simulations of further shock wave propagation in an adjacent tube filled with air or water (see section 5.3 applying an air-filled capillary).

4. Experimental setup

The experiments are performed with different commercially acquired tubes, namely glass capillaries (CM scientific) with inner widths of \( D = 50 \mu m, 100 \mu m, 200 \mu m \) and \( 300 \mu m \), respectively. The length of all tubes is \( L = 50 \) mm, the tube wall is half the thickness of the inner diameter of the tube.

The shock wave propagation in the tube is investigated by a laser differential interferometer (LDI), which is a modification of the arrangement published earlier in [50, 16]. The experimental setup is presented in figure 11.

The intensities of the two interferometer beams \( I_p(t) \) and \( I_s(t) \) are measured by two photodiodes (rise-time 10 ns), which yield the electric signals \( U_p(t) \) and \( U_s(t) \), correspondingly. The two photodiodes are integrated into an self-made electrical circuit, which stands for a compromise between the rise-time and the sensitivity in amplitude. The output signal of the circuit is the difference of the two photodiode signals \( U(t) = U_p(t) - U_s(t) \) as a function of time \( t \) or of the corresponding coordinate \( x \), respectively. \( U(t) \) is linearly proportional to the detected light intensity, which is modulated by the interference. The interference signal is a result of the phase difference (between the two interferometric beams), which itself results from the difference in the indices of refraction within the volume of the beams. Finally, the index of refraction correlates with the flow density via Gladstone–Dale relation and, therefore, the change in flow density caused by the shock wave can be deduced from \( U(t) \).

As shown in figure 12, the \( U(t) \) signal jump caused by the shock front is detected. Thus, the LDI is a suitable instrument to diagnose the density jump across the shock front \( \rho_2 / \rho_1 \). The shock velocity can be obtained through time-of-flight method by \( u_x = \Delta x / \Delta t \), thus \( u_x \) is averaged over the distance \( \Delta x \) between both interferometric beams and \( \Delta x \ll L \). Furthermore, the LDI yields the time \( t_i \) that the shock wave needs to propagate from the plasma to the first beam of the LDI at position \( x_i \), which determines the trajectory of the.
shock wave. As usual, applying the Rankine–Hugoniot relation, $\rho_2$ can be calculated from shock Mach number $M_1$ or velocity $u_s$.

Here it shall be remarked the following: this diagnostic method is different from most other diagnostics applied to shock tubes, which are predominantly based on pressure sensors (those often enable much worse resolution).

5. Results

5.1. Shock wave generation

Prior to the experiments of shock wave propagation in different tubes, various focusing sizes (all Gaussian-shaped) have been tested. Shock waves generated with the driver diameter $\phi < D$ are compared with those with $\phi \approx D$. At very early times $t \approx \tau_I \ll \tau_0$, the driver geometry should be planar in general for the present conditions (for fs-LPP, initially $d_{LPP}$ is only slightly larger than $d_{AI}$ and thus we have always $d_{LPP} \ll \phi$).

The driver geometry is point-like, if (I) $\phi$ is much smaller than $D$, or (II) the tube is absent.

Here it may be mentioned that despite the high laser power supplied to the glass plate, the laser pulse intensity is always below the threshold for self-focusing.

![Figure 11](image1.png)

**Figure 11.** Scheme of the experimental setup. The $1/e^2$-diameter of each interferometer beam inside the tube is 20 $\mu$m, the spacing between $\Delta x = 370$ $\mu$m. The distance $x_i$ from the LPP to the first interferometer beam is varied during experiments.

![Figure 12](image2.png)

**Figure 12.** A typical oscillograph trace of the shock wave detected by the LDI.
The focal lengths and f-numbers are given in Table 1 along with other parameters for the shock generation process.

Figure 13 shows the experimental results of shock wave attenuation affected by the propagation geometry, which is characterized in Table 1. As expected for fixed driver conditions, the shock wave is the strongest for the confined geometry of a tube \((D = 50 \, \mu m)\) driven by a quasi-planar LPP with a laser focal diameter \(f = 21 \, \mu m\). This diameter has been restricted to avoid additional excitation of the tube walls by the LPP. However, it is observed later on that such an effect is of minor importance. Even a larger focus does not change the result significantly as long as \(I\) is kept constant for all the investigated tubes.

On the other hand for a point-like driver \((f = 6 \, \mu m)\), the shock wave within the tube is strongly damped with increasing distance \(x\) from its onset position. This has been confirmed for the same intensity. But even for an order of magnitude higher intensity due to a smaller focal spot at constant laser energy, the shock becomes considerably weaker with increasing propagation distance (see Figure 13). Actually, if \(I > I_{\text{glass}}\), the effect of such an intensity variation on the LPP within the Al-layer is not much pronounced. However, no shock waves could be observed in case of the missing Al-layer.

Without the tube for geometric confinement, the shock wave in free air has a 3D propagation. The pressure behind the shock front decreases rapidly.

All this is consistent with a simple estimate based on the Sedov–Taylor expansion approach [14], where the volume expansion is modeled according to the corresponding geometry [16]: (i) accounting for the restrictions given by the cylindrical tube geometry and assuming a planar driver, (ii) supposing a spherical geometry in the early phase of shock evolution, i.e., for \(x < D\), which later on turns to a cylindrical geometry (additional reflections at the tube walls, see sections 2 and 5, may also contribute to shock wave attenuation, but this is not considered here) or (iii) treating a volume expansion without any constraints (when the tube is absent).

Like conventionally generated shock waves in tubes, the LPP-induced shock wave also has a speed-up phase \((\tau_b)\). But due to optomechanical limitations of the velocity measurements, the speed-up phase cannot be resolved in the experiments with tubes applied in the present work.

As a consequence of the geometric confinement, the planar shock wave in the 50 \(\mu m\) tube has a total propagation distance of approximately 2000 \(\mu m\) (then it becomes a sound wave). This rather long propagation distance of 40 times the tube diameter is of great interest for prospective investigations.

Due to these findings, all the following experiments with different tubes were restricted to planar geometry with focal spot diameters and intensities shown in Table 1.

### 5.2. Shock wave propagation in different tubes

As a first experimental investigation of LIMS propagation, experiments are performed for three different tube diameters, namely \(D = 50, 100,\) and \(200 \, \mu m\). Figure 14 shows the resulting positions of the shock front as a function of propagation time.
Obviously, the shock trajectories are not linear for all tubes, which corresponds to the shock wave attenuation (compare figure 13). Moreover, after a certain propagation distance, $x_{sw}$ (or the equivalent time $t_{sw}$), the wave is slowed down to sound velocity (index sw, a straight dashed line in the diagram).

Due to the high reproducibility of the measurements, the flow parameters can be obtained by fitting an allometric function $x(t) = a \cdot t^b$ that consists only of two parameters $a$ and $b$. This allows a direct comparison with the Sedov–Taylor similarity solution for a point-shaped explosion.

The flow velocity may be written in the form:

$$u_t(t) = \dot{x}(t) = abt^{b-1} = x(t) \frac{b}{t}.$$  

By setting $\dot{x}(t_{sw}) = u_a$ ($u_a$ is the sound speed), the critical values $t_{sw}$ and $x_{sw}$ can be obtained as

$$t_{sw} = \left(\frac{ab}{u_a}\right)^{\frac{1}{b-1}},$$

$$x_{sw} = a \cdot t_{sw}^b.$$  

Figure 15 shows the shock wave attenuation in different tubes. The two different methods—either via time-of-flight or via trajectory measurements—measure the same shock velocity within the 10% error range. However, the trajectory measurements provide higher spatial resolution, which is of particular interest for the early stage of the shock propagation (near field), since they have smaller steps between measurement positions (namely...
\( \Delta x_1 = 100 \, \mu m \) for the first 1000 \( \mu m \) of the propagation distance and \( \Delta x_1 = 200 \, \mu m \) for larger distances. The resolution of the time-of-flight method is determined by the step \( \Delta x = 370 \, \mu m \).

Figure 15(a) clearly shows that shock waves generated with the same laser intensity in bigger tubes propagate faster (are stronger) than those in smaller tubes. Here, velocities are plotted in absolute values.

This may be explained by the fact that a bigger tube offers less influence of friction and heat transfer than a smaller one. On the other hand, the shock propagation in small tubes quickly becomes laminar. In this context, it has been observed that the Reynolds number directly behind the shock front (i.e., downstream) in case of a 200 \( \mu m \) tube exceeds the value 2300 for at least \( t/t_{sw} = 0.4 \) and thus the propagation is turbulent. In contrast, the shock propagation in a 50 \( \mu m \) tube becomes already laminar for a Reynolds number larger than 0.1. Details of those investigations are beyond the scope of the present work and will be discussed elsewhere \[51\].

The observed attenuation in the initially turbulent motion may be explained by expansion waves reflected from the bottom of tiny tubes (refer to the Navier–Stokes computations in section 5.3).

However, the shock wave Mach numbers will not differ significantly, if the propagation distance is normalized to the tube diameter. (see figure 15(b)). This becomes evident, when values normalized to \( x_{sw} \) and \( t_{sw} \), respectively. Figure 16 shows that the corresponding normalized shock trajectories are almost the same.

From this we may state that much less laser energy (at fixed laser intensity) is needed to achieve the same shock strength (with respect to normalized values) in small tubes than in those with larger diameter (see table 1).

### 5.3. Comparison of Navier–Stokes computations with experimental results

Compressible laminar unsteady viscous flows in a micro tube are governed by unsteady axisymmetric compressible Navier–Stokes equations coupled with the multi-species conservation equations for a mixture. This set of equations may be written in a compact integral conservative form as

\[
\int_V \frac{\partial U}{\partial t} \, dV + \int_S F \, dS - \int_S G \, dS = 0,
\]

where the volume of a computational cell is denoted by \( V \) and its surface by \( S \). The definitions of \( U, F \) and \( G \) are:

\[
U = [\rho, \rho \vec{V}, E]^T,
\]

\[
F = [\rho_i (\vec{V} \cdot \vec{n}), \rho (\vec{V} \cdot \vec{n}) + p \vec{n}, (E + p)(\vec{V} \cdot \vec{n})]^T,
\]

\[
G = [\rho_i V_i^d, \vec{\tau} \times \vec{V} + \vec{q} \cdot \vec{n}]^T,
\]

with \( \vec{\tau} \) as shear stress and \( \vec{n} \) as unit vector; \( \vec{\tau} \) and \( \vec{q} \) denote the viscous stress tensor and the heat flux vector, respectively. Quantities \( \rho, p, \vec{V} = [u, v]^T \) and \( E \) are the density of the \( i \)-species, the pressure, the velocity vector and the total energy per unit volume, respectively; \( V_i^d \) being the diffusion velocity of the \( i \)-species. The subscript \( i = 1, 2 \) represents the species involved in the driven mixture.
Energy $E$ can be calculated by the relation

$$E = \rho \left( e + \frac{v^2}{2} \right),$$

(11)

where $e$ is the internal energy per unit mass defined as

$$e = \sum_{i=1}^{N} Y_i e_i(T),$$

(12)

with the mass fraction of each species, $Y_i = \rho_i/\rho$, and the density of the mixture $\rho$. The specific internal energy for each species may be expressed as

$$e_i(T) = \frac{3}{2}RT + \psi_i(e_{\text{tot},i}(T)),$$

(13)

with $\psi_i = 0$ for atoms or 1 for molecules.

Finally, pressure $p$ will be determined from the Dalton law: $p = \sum_i p_i$, where $p_i$ is the partial pressure of the $i$-species, assumed to behave as a perfect gas following the relation $p_i = \rho_i \frac{R}{M_i} T$. Here, $M_i$ is the mass per mole of the $i$-species, $R$ the universal perfect gas constant. Knowing the mass fraction of the species, their densities can be found from relations $\rho_i = \rho Y_i$.

The numerical solution of these equations was performed by using the parallel version of a multi-block finite-volume home code [18, 52] with an exact Riemann solver coupled with an AUSM-DV solver assuring second-order MUSCL extrapolation for the inviscid fluxes. The viscous and heat transfer terms were discretized using a central difference scheme. Grid cells were refined near the wall with a minimum non-dimensional $y/D_{\text{H}}$ step equal to $10^{-2}$ at the wall, and the mesh size is 1000 in $x, y$ directions. The resulting integration time step was $10^{-11}$ s.

As in experiments, three hydraulic diameters $D_{\text{H}}$ are chosen equal to 50, 100 and 200 $\mu$m. The driver tube length is 2.5 $\mu$m, in which the discharge is simulated by a given value of the pressure $p_{\text{1}}$ at $t = 0$. The initial temperature in the tube is equal to 300 K and the driven pressure $p_{\text{1}}$ is atmospheric one.

On the solid walls, the following boundary conditions were used:

$$u = v = 0, \quad T = T_w, \quad \frac{\partial p}{\partial n} = 0,$$

where the subscript $w$ refers to the wall quantities.

The first test case in a 200 $\mu$m diameter tube has been computed with an initial pressure ratio $p_{\text{1}}/p_{\text{1}} = 100$. Although this pressure is rather low when compared to the corresponding $p$ deduced from the MULTI-fs hydrocode simulations (section 3), it must be stated that $p_{\text{1}}/p_{\text{1}} = 100$ indicates the ‘effective’ pressure ratio for shock formation. The pressure distribution along the center line of the tube at different times is plotted in figure 17. The computations stop when the shock wave reaches the tube exit ($t = 10^{-3}$ s).

It can be clearly seen that the shock wave attenuates during its propagation along the tube with decreasing pressure peaks. One can also see the wave expansion and its reflection at the bottom of the tube. A secondary shock wave appears, which brings the pressure back to the atmospheric one. Between the main and secondary shock waves, the pressure decreases and thus this behavior is different to the classical shock tube, where a plateau pressure appears behind the shock wave.

In order to show the flow structure in the tube, the axial velocity contours are drawn at three different times in figure 18.

The main and secondary shock propagation and the wall boundary layer development between them are clearly visible. The distance between these two shock waves increases with time. The intensity of the shock wave velocity along the tube can be deduced from the computations and is plotted in figure 19. The numerical results agree well with the experimental data and tend to validate this numerical description.

The same computations have been conducted for the two other diameters 100 and 50 $\mu$m with the same mesh size. The main difference is the choice of the initial pressure ratio which has been reduced to 65 for these two cases, in order to have the best fitting with experimental data as shown in figure 20, where the influence of the initial pressure ratio is tested on the shock velocity evolution along the 100 $\mu$m diameter tube.

The main requirement for the appropriate choice of the initial pressure ratio is to obtain a good agreement with experimental data on the first part of the tube (approximately 20 diameters). It can be also noted that whatever this ratio is, the shock velocity tends to the same final value (sonic value).

For the investigated cases with different tube diameters, the shock wave velocity along the tube is plotted in figure 21, which shows good agreement of the numerical values with experimental data. The slope of attenuation along the tube and the limiting values are also well described.
The observed behavior of the shock wave (main) velocity decline along the tube can be explained by two processes. The first one is the core flow expansion behind the created shock wave which is caused by the initial pressure ratio and a very short driver. The second one is the development of the wall boundary layer which interacts with the core flow. These two processes lead to a decrease of the shock wave velocity.
Figure 22 shows the influence of each process on the shock wave attenuation. The main effect is the first one. It is independent on viscous flow effects as obtained from Euler computations (labeled Euler in figure 22). The Navier–Stokes computations (labeled NS) with an open entrance tube and the same pressure ratio allow also to capture the wall boundary layer development behind the shock wave, which additionally reinforces its attenuation (labeled BL). This latter numerical result agrees well with the correlation: shock wave velocity versus local scaling factor $Sc_x$ (labeled correlation), which is given in [20] and obtained in classical micro shock tubes.

The shock wave velocity evolutions gained from different model calculations are compared to the experimental data for a 100 $\mu$m tube (see figure 22). A similar comparison can be observed for the two other diameters (not shown here).

To sum up, the propagation behavior of shock waves in the investigated micro tubes, where the shock wave is initially produced in a laser plasma, is mainly determined by the expansion of the flow behind the shock front along the tube. This leads to a decrease of the shock wave intensity. The effect is compounded by the influence of the wall boundary layer, which interacts with the core flow and strengthens the decay of the shock wave. The whole propagation scenario is well described by solving the unsteady laminar compressible Navier–Stokes equations. This approach is confirmed by the agreement between experimental and numerical findings.
6. Summary and conclusion

In summary, a novel method to generate shock waves in small tubes with diameters down to the micrometer range has been developed, namely LIMS, which makes use of plasmas generated in thin layered targets by means of intense femtosecond laser pulses. To the best of our knowledge, the present work is the first report on direct micro shock wave generation and propagation in currently available smallest diameter tubes at clearly defined conditions. Thus, it opens access to the experimental shock wave research on micro scale and hence strongly contributes to this new topic that was mostly restricted to theoretical work so far.

It has been shown that by a careful design of the LPP as the driver, strong shock wave generation in air is possible. Different scenarios of this process have been investigated. In particular, shock wave generation in free air and under confined conditions of small tubes, both in point-like and in planar driver geometry, was tested. As expected, planar laser plasmas are the most suitable drivers and offer the opportunity for the propagation of strong shock waves in tubes with sub-mm or micrometer diameter. In such a confined geometry, propagation lengths up to several mm become available, before the wave converts to a sound wave. Experimental results have been supported by hydrocode simulations. Both show good agreement.

As a first application of LIMS, micro shock wave propagation has been studied in tubes of different diameters between 50 and 200 μm. Optical diagnostics, in particular, a LDI, with the advantage of a non-contact
measurement, were applied to investigate the shock wave behavior, in particular, shock wave attenuation. Strong shock waves could be observed for all applied tube diameters. Moreover, applying the same laser intensity in the experiments, all shocks exhibit the same numerical properties, as far as the propagation distance (or time) is normalized to the distance (or time), at which the shock wave velocity approaches the sound speed. Therefore, it can be stated that down-scaling does not lead to a weaker shock.

In addition, the observations clearly show that for the same shock strength, smaller tubes need less input laser energy, when compared to larger ones. Especially interesting and useful for applications is the fact that shocks propagating in smaller tubes do more quickly transfer from turbulent to laminar flow.

The Navier–Stokes computations allow a well description of the shock wave attenuation along these microtubes. They show that the attenuation is mainly caused by the core flow expansion behind the shock wave. Differently from classical shock tube physics, the boundary layer contributes for the less part to the shock attenuation here.

It may be concluded that the present work yields prospects for further fundamental research and may also trigger future applications, e.g., in medical physics or micro surgery.

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