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by

HEINRICH HEINE
UNIVERSITÄT DÜSSELDORF

Study of Space Radiation Effects with Laser-Plasma-Accelerators

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Study of Space Radiation Effects with Laser-Plasma-Accelerators

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B. Hidding\textsuperscript{1,2,4}, O. Karger\textsuperscript{2}, T. Königstein\textsuperscript{3},
G. Pretzler\textsuperscript{3}, J.B. Rosenzweig\textsuperscript{4}

\textsuperscript{1} Scottish Center for the Application of Plasma Accelerators, University of Strathclyde, Physics Department, Glasgow G4 0NG
\textsuperscript{2} Institut Für Experimentalphysik, Universität Hamburg, 22607 Hamburg
\textsuperscript{3} Institute of Laser and Plasma Physics, Heinrich-Heine-University Düsseldorf, 40225 Düsseldorf
\textsuperscript{4} Department of Physics and Astronomy, University of California, Los Angeles

Hamburg, Germany

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1 Executive Summary

The potential of laser-plasma-accelerators for space radiation testing of electronics is explored. Both fields, laser-plasma-acceleration on the one hand and space radiation testing on the other are highly vibrant fields, which have been disjunct so far. However, there are a multitude of aspects which suggests that laser-plasma-accelerators could be highly relevant future radiation sources for the space radiation testing industry. It is the goal of this study to connect both fields and to introduce laser-plasma-accelerators as complementary radiation sources for improved space radiation testing.

First, a brief review of state-of-the-art testing techniques in the space radiation testing community is given. There is a need for novel additional radiation sources, mainly because of the limited availability of beam time, because of the limited congruence of spectral flux in space and what can be produced with state-of-the-art sources, because of the flux limitations in conventional sources, and because of the limited flexibility of today’s space radiation testing facilities.

Then, the fundamentals and the status of laser-plasma accelerator technology are briefly reviewed. The main motivation for research in this rapidly evolving field stems from the fact that laser-driven plasmas can sustain electric fields of tens of GV/m, which can be used to accelerate particles such as electrons, protons and ions in a broad parameter regime in very compact setups. High-power laser systems are needed, which are today commercially available in various types. Currently, the best choice as regards particle acceleration performance are table-top sized Ti:Sapph solid-state lasers, although even more efficient and cost-effective laser systems such as thin disc and fiber lasers have shown dramatic advances over the past years and are expected to reach the pulse power levels required for efficient particle acceleration in the next few years. The increasing availability of high-power laser systems suitable for radiation generation in Europe is discussed, and the physical principles behind generating high-power laser pulses are summarized. An overview on the interaction of such laser pulses with matter (both overdense and underdense) is given, which includes laser energy absorption and reflection, ionization and plasma formation, electron and proton acceleration, propagation through the target, and particle radiation emission. Next to their compactness, outstanding features of laser-plasma-accelerators are their flexibility and the ability to generate important types of radiation which have been hitherto inaccessible with conventional radiation sources. Most types of radiation are very broadband, often spanning several orders of magnitude of energy. While conventional accelerators typically generate radiation with very narrow energy spread, production of broadband radiation is the natural domain of laser-plasma-accelerators. It is shown that this broadband radiation can be either tweaked towards monoenergetic beams, but can also be tuned to reproduce specific irradiation scenarios in space with high accuracy as regards spectral flux. Examples of state-of-the-art laser-plasma-acceleration experiments are given and important key diagnostics which are typically used in this community are presented.
Finally, by combining techniques and principles of both the space radiation testing community as well as the laser plasma acceleration community, which was possible in the context of the agency’s Networking Partnering Initiative (NPI) by having a joint research team from both communities, first proof-of-concept experiments were carried through. While laser-plasma-accelerators are capable to produce protons, ions, electrons and hard photons alike, careful design considerations in the preparation phase led us to concentrate mainly on a specific experimental scenario, namely the reproduction of electron radiation as present in Earth’s radiation belt. These electrons are often also called ”killer electrons” since they represent a strong hazard to space vessels operating inside the van-Allen belts, which includes the important geostationary orbits (GEO) such as navigation satellites. At the same time, laser-plasma-acceleration of electrons in the corresponding energy range is especially maturely developed already, which strongly suggested it as a proof-of-concept case. The design considerations indicated that laser-solid interaction was better suited than laser-gaseous interaction for this scenario, which we realized by using a strongly focused laser pulse incident on thin (µm-scale) metal foils. The focused laser intensity is many orders of magnitude higher than the ionization threshold of the foil, which is why the foil is turned by the laser into a plasma quasi-instantaneously. The rapid motion of the released plasma electrons in the ultraintense electromagnetic laser field leads to a rapid acceleration of electrons, mainly into the forward direction, but due to additional complex interaction mechanisms, also in specific other directions such as parallel to the target foil. The electron energy was tuned by varying the laser intensity on target to match the energy distribution of van-Allen belt electron flux in space. Due to the high electron density present in the solid density target foil, and the emitted electron flux duration being of the order of the incident laser pulse which is of the order of few tens of femtoseconds, and the spatial source size – the laser focal spot size – being as low as few micrometers squared, the emitted electron flux is very high initially after emission from the target. On the one hand, this constitutes a possibility to carry through extremely high peak flux studies, orders of magnitude larger than with conventional sources. On the other hand, because purposely very broadband radiation is generated, because of a huge emission angle of tens of degrees, and because of an intensity-dependent scaling of the generated beam temperature and divergence, the flux level can also be reduced seamlessly towards almost arbitrarily low values simply by tuning the distance of the target foil from the focus.

In benchmarking campaigns, the generated flux of electrons in the laser-plasma-lab was analyzed in great detail, and optocouplers were benchmarked by ESA in state-of-the-art space radiation test facilities. These optocouplers and other test devices were then irradiated in the laser-plasma-lab under well-known conditions in multiple campaigns and afterwards analyzed at ESTEC. To our knowledge, this was the first time that space radiation hardness studies were carried through with laser-plasma-accelerators. It was also the first time that killer electron radiation such as present in Earth’s radiation belt was reproduced in the laboratory. In these proof-of-concept irradiation campaigns,
dose-dependent damage was inflicted successfully on the devices under test. This shows for the first time that laser-plasma-accelerators are workable novel complementary advanced radiation sources for space radiation studies.
2 Introduction

Radiation in space is one of the major threats to manned and unmanned missions. With an ever increasing number and complexity of space missions, and at the same time increasing demands on the performance of electronics onboard space vessels, this fundamental problem is continuing to grow more and more important. European Space Agency and space entities all over the world are constantly developing and using various strategies and countermeasures in order to respond to this threat.

![Figure 1: Aurora Australis. Left: UV image taken by the Imager for Magnetopause-to-Aurora Global Exploration (IMAGE) spacecraft (2011). Right: Image taken from the ISS (2011). Figure credits: NASA.](image)

The paramount importance of radiation effects for space exploration fuels various research fields. First, one has to know what kind of radiation and what radiation intensity to expect during a mission specific mission. Experimental knowledge on that comes from observations and measurements. The use of Earth-bound observation is limited due to the magnetic field and atmospheric protection. Figure 1 shows the particle bombardment of Earth vividly via UV Earth observation images from space. At higher altitudes, the radiation level is significantly higher than on the Earth’s surface. This increases the radiation exposure dose onboard civil airplanes during intercontinental flights, for example. Since the 1930s, high-altitude balloons equipped with various diagnostics were used to measure the ionization effects of cosmic radiation and other radiation characteristics [10]. With the advent of the space age, the high radiation level in space was becoming ever more clear. For example, Explorer I, the first American satellite, discovered the van Allen belt in 1958. Solar activity and plasma storms can substantially distort the Earth’s magnetic field and is therefore one of the key influencer for space weather. Numerous missions have contributed since the 1960s to collect further data on radiation belts and space weather in general. Well-known past and current ESA/NASA/AF/JAXA missions include CRRES (Combined Release and Radiation Effects Satellite, 1990), FAST (Fast Auroral SnapshoT, 1996), SAMPEX (Solar Anomalous and Magnetospheric Particle Explorer, 1992), GEOTAIL (1992), IMAGE (Imager for Magnetopause-to-Aurora Global...
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Exploration, 2000), SOHO (Solar and Heliospheric Observatory), Cluster (studying the Earth’s magnetosphere over an entire solar cycle), Double Star, THEMIS (Time History of Events and Macroscale Interactions during Substorms, 2007) and ARTEMIS (Acceleration, Reconnection, Turbulence and Electrodynamics of the Moon’s INteraction with the Sun), and CINDI (Coupled Ion Neutral Dynamic Investigation, 2008). More recently, the Radiation Belt Storm Probes (RBSP) had their lift-off in 2012, and are supposed to bring detailed further knowledge on particle acceleration in the radiation belts, and the influence of the sun and geomagnetic storms on the radiation belts. For example, they already have discovered a transient 3rd van-Allen belt in 2013 [11].

These experimental approaches are complemented by theoretical analysis and computational modeling, which is a very active research field in its own right. The injection, trapping and acceleration processes in the radiation belts, for example, are extremely multi-faceted, involving aspects as diverse as solar activity and flares, geophysics and electromagnetic and plasma theory.

Radiation in space can be extremely versatile: There is a plentitude of different types and origins of space radiation. The mechanisms responsible for the generation can be highly complex and at the same time very different from each other. However, no matter whether it is particle radiation consisting of electrons, protons, ions or photons, there is one attribute which holds for all known kinds of space radiation: the radiation is never monoenergetic. Instead, the radiation is always very broadband, often spanning multiple orders of magnitude as regards energy. This is in diametral contrast to the typical radiation sources used in today’s space radiation testing facilities, where often cyclotrons and electron linacs are used, which produce highly monoenergetic radiation. Furthermore, while a space vessel mostly encounters multiple different kinds of space radiation, for example electrons, protons and ions, a linac or cyclotron is typically capable of producing only one specific kind of particle flux. This often necessitates multiple, time-consuming and costly test runs at different facilities. Flexible radiation sources, with the capability to produce broadband radiation of different types, ideally with high flux in order to keep reduce testing times are therefore desirable.

In addition to the complexity of the space radiation environment, solar activity heavily influences the incident radiation on wide space- and time-scales. For example, eruptions in the surface of the sun (Coronal Mass Ejections, CME) can generate shocks in the solar wind, which heavily distort the geomagnetic structure. This can change the radiation level on the timescale of hours or even minutes, and the associated phenomena are therefore called space weather. Prediction of space weather is therefore a very important, multi-faceted goal in the context of space radiation. The prediction of space weather is complicated by the distortion and compression of the Earth’s magnetic field lines by the solar wind on the sun-facing side, and the expansion on the opposite side. Today, various websites run by international organizations such as ESA and NASA exist which display the daily space weather in a similar manner as the common public atmospheric weather is predicted.
Based on the more or less detailed knowledge of space radiation environment, testing of the sensitivity of electronics can be performed. The ideal situation were if one could reproduce the radiation as expected in space during certain stages of the mission one-to-one. In practice, one is limited to using radiation sources which are available here on Earth. The radiation output from these sources mostly is only a very rough approximation of what is really occurring in space. Therefore, the accuracy and significance of testing procedures is limited by

1. the knowledge and predictability of the radiation composition during various stages of the mission

2. the available radiation sources used to reproduce the expected space radiation

In addition, since space radiation consists of electrons, protons and ions, and since space radiation can have various different effects on electronics, realistic reproduction of space radiation requires several testing procedures to be carried out. Because radiation sources are typically not flexible enough to produce protons as well as electrons in a wide parameter range as regards energy and flux, different facilities have to be used to investigate the different effects on electronics. For example, a device is tested not only at a cyclotron which produces protons and/or ions, but also at a linac or electrostatic accelerator which produces electrons. Furthermore, the sometimes large fluence which occurs in space during a mission would require to irradiate the device under test (DUT) over an extended period of time – which can be several years, which often places experimentally prohibitive hurdles – in order to accumulate the same fluence as in space. Since beam time is a valuable commodity, these factors add up to complex, costly and time-consuming testing. In this scenario, various tradeoffs have to be made between cost, duration and significance of testing.

Based on the radiation and testing requirements, the selection process of electronic devices for specific missions is carried out. The devices to be used range from Commercial-Of-The-Shelf (COTS) electronics, for example for missions and satellites operating at Low Earth Orbit (LEO) or in the "safe zone" between the inner and outer radiation belt, where radiation levels are relatively low, to radiation-hardened-by-design (RHBD) devices for extreme radiation environments such as Jovian orbits. Device selection and development goes hand in hand with determination and design of shielding in order to decrease the flux levels incident on the electronic systems onboard. Hardness assurance must be addressed at various levels, for example at the parts, board, package and system level. This further increases the need for beamtime and sophisticated testing methods.

Increasing the accuracy of significance of testing procedures can lead to substantial increase of predictability of the behavior of electronic parts during missions, and the mission reliability. Further, increasing the availability of beamtime for radiation hardness tests is also one major necessity for space exploration. The number of space-exploring...
countries and missions increases, and at the same time the performance and size of electronic devices decreases, which often increases the radiation sensitivity of these devices. Both add up to an ever increasing demand for beamtime. Additional facilities are therefore needed to provide new radiation sources capable of providing usable beamtime, and advanced testing methods have to be developed.

To sum up, the various research areas which are building blocks for mission planning are

- Experimental space radiation measurements
- Theoretical modeling of space radiation
- Space weather prediction
- Testing of radiation effects on electronics and biological systems following established procedures
- Development of rad-hard electronics and countermeasures
- Development of advanced testing procedures and radiation sources

The activity which has led to this final report aims mainly at the last one of these building blocks, which might be regarded as the most fundamental one. The development of advanced testing procedures and radiation sources has the potential to increase fundamentally the significance of testing methods, which in turn would lead to enhanced predictability in mission planning, and ultimately, to increased reliability of missions.

3 Space Radiation Testing Techniques Today

Radiation in space is extremely multi-faceted. Important main components of space radiation are

- radiation belt electrons and protons
- solar proton events (SPE)
- galactic cosmic rays (GCR)

These different types of space radiation, often being mixed and/or occurring during different phases of the mission, require multiple facilities to be involved in space radiation effects testing. The problem is that hitherto there are no facilities which would be versatile enough to provide test beams for all of the mentioned types of space radiation. What’s more, there are hitherto few dedicated facilities which can provide higher energy electrons, protons and ions to the user. Rather, space radiation testing often has to
compete with other R&D campaigns at research radiation sources. In general, one can say that there is a chronic shortness of beam time at these facilities, and allocation of beam time is typically a costly and time-consuming process.

Electrons, protons and ions can occur in extremely broad energy range, from keV to the most energetic particles with energies \(> 10^{21}\) eV. As pointed out in [1], for example, the majority of space missions are taking place primarily in Earth orbits, for example at geosynchronous orbits. It is just here, where energetic populations of electrons and protons are located in the van-Allen radiation belts, and therefore are especially dangerous and important. Figure 2 shows an artist’s view of the general radiation belts’ geometry.

![Figure 2: Artist’s view on Earth’s van Allen belts (image courtesy of NASA).](image)

The major component of the total dose on a typical orbit is contributed by the electron flux, a large fraction of which is coming from the electron flux trapped in the radiation belts. This is shown in figure 3. Here, the yearly dose rate after about 1 mm of aluminum spacecraft shielding for a low Earth orbit with 60° inclination is plotted. As can be seen (note the logarithmic scaling of the y-axis), the ionizing dose to be expected due to electrons is an order of magnitude higher than the contribution due to protons up to an altitude of about 1000 km, and is still substantially higher than the proton contribution at a distance of about two Earth radii. This underlines the high importance of radiation belt electrons for space radiation damage.

Detailed knowledge, understanding and predictability of space radiation is desirable in order to assess in a reliable way the performance of electronics during the actual mission. Due to the complexity of the radiation environment during the whole mission lifetime, the most realistic and straightforward method for testing these electronics would be to test them in the actual space environment. Such missions have actually been flown. One
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Figure 3: Radiation belt electrons can often be the dominating contributor for the total dose experienced on an Earth orbit (from [1]).
example is the Combined Release and Radiation Effects Satellite (CRRES) in 1990/1991. However, such a direct approach has a number of disadvantages. The most obvious one is the enormous costs connected to such an approach. But also the very long time scales required, add up to making this approach practically prohibitive. As pointed out in [1] and [12], for example, a desirable alternative is "to duplicate the space environment to the greatest extent possible" in ground-based laboratories [1]. However, this can be challenging due to the complexity of the space radiation. Various parameter regimes and radiation sources may be necessary in order to reproduce space radiation. One feature of space radiation is that they typically occur not monoenergetic, but instead in a broad energy band. Using traditional radiation sources, such broadband flux is very difficult to realize experimentally. Therefore, following [12], radiation testing is mostly done "under the following simplifications: usage of monoenergetic electron or ion beams instead of fluxes of particles with distributed energy spectra", which requires "serious scientific ground, i.e. detailed knowledge of physical mechanisms of the radiation effects to avoid the wrong results."

It is one of the main goals of the present ESA NPI activity to show that exact reproduction of broadband space radiation in various regimes is possible by introducing laser-plasma-accelerators as novel radiation sources for space radiation testing. This, as well as the necessity to produce broadband, exponential electron flux for more realistic testing, and the difficult using established technology, was pointed out in [13]. In the following, established laboratory radiation sources and testing techniques will be briefly reviewed along these lines.

Today, various types of charged particle accelerators are commonly used for simulation of charged particle space radiation such as electrons, protons and ions. Among these are electrostatic devices such as the Dynamitron as well as linear accelerators, cyclotrons and the like. Smaller, compact radiation sources such as cobalt-gamma sources can be maintained as dedicated facilities, but generally for higher particle energies, the size, maintenance and costs of state-of-the-art accelerators are prohibitive to have them on site of the company/entity which develops the electronics and/or tests its radiation hardness.

One example of a rather compact device is the Dynamitron. It has been invented in 1960 [14], and has then made quite large industrial impact being distributed by Radiation Dynamics, Inc. (compare figure 4, left hand side). For example, one large Dynamitron is installed at the Jet Propulsion Laboratory, where it is used for space radiation testing. It can deliver up to 2 MeV electrons as well as protons by acceleration on a static electric field. The whole evacuated device, as pictured in the right hand side of figure 4 has a diameter of approximately 2 m.

In 1999, Radiation Dynamics, Inc., was purchased by Ion Beam Applications, a company which originally was a spin-off of the Cyclotron Research Center of the Catholic University of Louvain-la-Neuve (UCL) in Belgium. The cyclotrons at UCL are today still heavily used by European space entities as well as by ESA for space radiation test-
Figure 4: Large (approx. 2 m diameter) Dynamitron installed at Jet Propulsion Laboratory, Pasadena, USA. Photo taken with permission by B. Hidding at JPL in 2011.
ing, and were used for testing of components relevant to the present ESA NPI activity as well as for diagnostics calibration during this project. Figure ?? shows the worldwide distribution of Dynamitron devices.

ESTEC has dedicated facilities on site and others, partially as part of the European Components Irradiation Facilities (ECIF). There is a 2000 Ci-level $^{60}$Co $\gamma$-source on site at ESTEC’s ECF (Electronic component facility). Here, many components can be tested at the same time in a huge room, with variable flux and fluence by varying the distance to source and/or the total irradiation time. While there is no direct control on the radiation on the radioactive Co source, the irradiation on the samples can be switched off and on within few minutes by driving the source in and out of a massive radiation shield. Such a source is elegant due to its simplicity, but on the other hand it is well-suited only for total dose tests, and has obvious drawbacks connected with handling of radioactive substances, including environmental and proliferation aspects.

Same holds for CASE, the Californium-252 Assessment of Single-event Effects laboratory at ESTEC, loaded with a few mCi Cf-252 source. Next to these ESTEC-internal facilities, there are external facilities which are used by ESTEC and partners on a regular basis, many of those as part of the ECIF, the European Components Irradiation Facilities. Among these are

- proton and ion sources at UCL at Louvain-la-Neuve
- RADEF in Finland
PIF at PSI, Switzerland

and other linacs and cyclotrons, subject to available beamtime.

Radiation in space is generated by stochastic acceleration mechanisms. These mechanisms naturally lead to very broadband spectra, which is summarized in [15]. In more detail, the radiation is not only broadband, but characterized by an exponential decrease of spectral flux towards higher energies. Again, this is an effect of the underlying acceleration mechanisms: it is less probable that a particle is accelerated to higher energies than to lower energies. This general law holds for electrons, protons, ions, and neutrons alike. In [15], figure 1, this is shown for generalized cosmic protons, in figure 2 for neutrons, in figure 3 for Earth Radiation Belt (ERB) protons as well as GCR protons and neutrons inside the Columbus module, in figure 6 for Jupiter radiation belt electrons and protons etc. All these spectra are characterized by the broadband spectral flux, which is strongly decreasing towards higher energies.

Now in sharp contrast to that, today’s space radiation test facilities and procedures are characterized by that the linacs and cyclotrons used do not produce broadband spectral flux at all, but instead strictly monoenergetic beams. Following [12], this is clearly a dramatic simplification of the realistic scenario.

Figure 6 illustrate the difference in stopping of a monoenergetic electron beam with \( E = 4 \text{ MeV} \) and an exponentially distributed one, with an effective temperature of 2 MeV after passing a shielding. The protection shielding consists of a multi-layer of Aluminium, each 1 mm in thickness. The marks (i) - (iv) denote the flux after 1, 2, 3, and 4 mm aluminum. The monoenergetic beam decreases in maximum energy and is also getting more broadband (fig. 6a), and in diametral contrast to the space scenario there is a lower flux at low energies than at high energies. In contrast to the monoenergetic beam, the exponential-energy beam remains exponential with the exception of a low-energy cutoff (fig. 6b). The ionizing dose shows a similar behaviour (fig 6c and d). The TID increases during the passage through solid aluminum in case of the 4 MeV beam and reaches its maximum after about 3 mm, when a large fraction of particles has been decelerated down to low energies and their energy is deposited. In contrast, with the exponential energy flux the TID decreases constantly during the passage through matter. Therefore the clear conclusion can be drawn that broadband incoming radiation with a flux as similar as possible to the one in space would be much better suited for space radiation than using monoenergetic beams such as obtainable from linacs and cyclotrons today. This is the natural regime of laser-plasma-accelerators! As we will see in the following sections, in laser-plasma-accelerators similar radiation producing mechanisms are at work as in space environments. A consequence, the radiation produced by laser-plasma-accelerators can be relatively straightforwardly tuned to reproduce space radiation here in Earth in the laboratory.

In addition to the highly beneficial spectral flux, laser-plasma-accelerators have at least two more fundamental advantages when compared to traditional radiation sources. First,
Figure 6: Comparison of spectral flux change of (a) an incident monoenergetic electron beam (as from a linac) of $E = 4\text{ MeV}$ when straggling through matter, and (b) an exponential beam as from an LPA with $T_{\text{eff}} = 2\text{ MeV}$. The marks (i) to (iv) denote the forward flux after passing through 1 to 4 mm of aluminum. In contrast to the monoenergetic beam, the exponential-energy beam remains exponential with the exception of the low-energy cutoff. (c) and (d) show the total ionizing dose (TID), which is deposited in a 6 mm thick aluminum wall after irradiation by monoenergetic and exponential beams, respectively.
they are highly versatile and can produce electrons, protons and ions in a very similar experimental setup, and can even produce electron and proton-ion flux at the same time – just as what is the case in space. Second, the peak flux of laser-plasma-accelerators can be tuned to enormous values, which are by far higher than what is present in space, such that this can potentially be used not only to decrease the irradiation times – e.g. for testing of the degradation of electronics over a whole satellite life – but also to test highly nonlinear regimes, where additional radiation damage effects can be studied.

In the following sections, in order to provide background information on the principles (the description of this comes in large parts from refs. [9], previously unpublished in English) and the capabilities of laser-plasma-accelerators, details on the history and physical and technical foundations of laser-plasma-accelerators as well as current and future trends are given. It is shown that space radiation reproduction is a self-evident application of laser-plasma-accelerators with potentially transformative impact on the radiation testing community. To develop this application to a complementary and highly synergistic work field is one of the main recommendations of this project report.
4 Review of Laser-Plasma-Technology

In this section, the development of laser-plasma accelerator technology is summarized. Due to various technological breakthroughs as regards laser pulse generation, LPA technology is today available commercially throughout the world. Costs and reliability of laser systems are steadily decreasing, while the performance and reliability of these systems steadily increases. In turn, this results in ever increasing distribution of laser systems and market penetration. Today, a large variety of laser systems is in widespread use both in research and industry. Here, for reasons to be made clear later, we seek for compact laser systems with high pulsed power, high reliability, good affordability, easy maintenance and with environmentally-friendly footprint.

Rutherford and his team demonstrated some 100 years ago in seminal experiments, that atoms consist of massive positively charged nuclei, surrounded by negatively charged electrons with much lower mass. This discovery contributed fundamentally to the advent of accelerator physics. While Rutherford used radioactive substances to generate the particle radiation which was necessary to come to his conclusion, later more effective particle sources and accelerator devices were developed, such as linear accelerators (linacs) and cyclotrons. Later, at a conference at CERN in the 1950s it was suggested for the first time to use the collective fields for acceleration which arise when the very building blocks of matter as discovered by Rutherford, namely electrons and ions, are temporarily separated from each other [16–18]. Since fully ionized matter is called a plasma, this strategy later became known as plasma acceleration. Today, plasma acceleration is seen by many as one of the most advanced particle acceleration method for various reasons. Perhaps the most fundamental one is that the accelerating fields in plasma accelerators can be many orders of magnitude larger than in traditional accelerators.

Most generally speaking, the energy a charged particle can gain in an accelerator is given as \( W = qEd \), where \( q \) is the charge, and \( E \) is the averaged electric field the particle experiences over the acceleration distance \( d \). Mainly because the electric fields \( E \) in conventional accelerators are limited to few tens of MV/m, the demand for higher and higher particle energies could be satisfied for the most part only by introducing ever longer acceleration distances \( d \). Prominent examples of such large accelerator facilities such as SLAC, HERA, and the LHC are depicted in figure 7.

The collective fields which arise [16–18] when provided by transient electron-ion separation in longitudinal plasma waves [19, 20] can be about four orders of magnitude stronger than the maximum fields in conventional accelerator technology and thus, in principle, allow similarly powerful accelerator constructs on the sub-m-scale.

Figure 8 visualizes the maximum accelerating electric fields obtainable in conventional radiofrequency-based accelerator modules in the L, S,X,V, and W band, contrasted by the maximum electric fields in plasma via the particle beam driven plasma wakefield acceleration (PWFA), laser wakefield acceleration (LWFA) and the self-modulated laser
wakefield acceleration (SMLWFA) schemes. The electric fields in focused laser pulses can be even higher than the peak fields in plasma, but they are not easily directly usable for particle acceleration to high energies, since they are oscillating transversally. Generally speaking, one needs charge separation in a plasma in order to convert the transversally oscillating laser pulse field into a longitudinal electric field suitable for acceleration.

In order to ionize matter and generate plasma, and separate at least for a short period of time the electrons from the ions, large electric fields are necessary. These can be provided either by dense particle bunches or by focused high-power laser pulses. Focused laser pulses can have (oscillating) electric fields which can be many orders of magnitude higher than the ionization thresholds of any matter, and also orders of magnitude higher than the electric fields associated with even the finest particle bunches available today. They are therefore ideal candidates in order to ionize, and to dispel electrons away from the positively charged ions. The electric fields $E$ generated by charge separation in plasma scale with the plasma electron density $n_e = m_e \omega_L^2 / (4\pi e^2)$ and is therefore tunable. Here, $m_e$ is the electron mass, $\omega_L^2$ is the laser angular frequency, and $e$ is the elementary charge.

There are several ways how to generate plasmas which can be used for particle acceleration. When using laser pulse drivers, a fundamental distinguishing mark is whether the laser-matter interaction is overdense or underdense. If the matter is overdense with respect to the driving laser pulse, the electron density in the matter is so high that the charge displacement due to electron movement is so effective that the propagating electric field will be reflected, whereas the laser pulse would be able to propagate through the ionized matter if the plasma is underdense. In experimental scenarios, there is always a more or less pronounced prepulse preceding the main laser pulse, for example due to amplified spontaneous emission (ASE) processes in the laser medium. Since prepulses can generate a preplasma, it is generally aimed at to suppress the prepulse level as far as
Figure 8: Comparison of electric fields in focused laser systems and beam- and laser-driven plasma wakefields (from [21]).

possible. In case of initially overdense matter, the point where the (pre)plasma changes from being underdense to being overdense is called the critical cutoff density $n_c$. As the density increases, the phase velocity of the electromagnetic wave increases and the group velocity increases. At critical density, the plasma frequency is equal to the light frequency and the light wave is reflected. The equation for the critical density is

$$n_c = \frac{\omega_L^2 \varepsilon_0 \gamma m_e}{e^2}$$

where the electron energy $\gamma$ takes into account relativistic effects. For Ti:sapphire laser systems, which have a central laser wavelength of $\lambda_L \approx 800$ nm, corresponding to an angular frequency of $\omega_L \approx 2.4 \times 10^{15}$ s$^{-1}$, the critical density amounts to $n_c = 1.7 \times 10^{21}$ cm$^{-3}$. Matter in the solid state, such as metals, typically has a typical electron density of the order of $n_e \approx 10^{22}$ cm$^{-3}$, corresponding to a plasma frequency of $\omega_p \approx 5 \times 10^{15}$ s$^{-1}$ or a plasma wavelength of $\lambda_p \approx 300$ nm. Therefore, metals are blocking visible light, while they might be transparent for UV light and even higher frequencies.

Conventional radiofrequency-based accelerator technology is fundamentally limited by material breakdown. This limits the accelerating fields in conventional accelerators. Even when using state-of-the-art superconducting magnets based on niobium-titanium NbTi or triniobium-tin (Nb$_3$Sn), the maximum accelerating fields when using state-of-the-art techniques are still limited to tens of MV/m. That is one of the core reasons why the increase of maximum energy gain in conventional accelerators has started to saturate.
in the last decades. Increasing the maximum energy would require bigger accelerators, which in turns requires resources which begin to exceed the justifiable level. In contrast, the relatively new plasma-based approaches, offer much more room for improvement. In classical accelerator technology, it is common to demonstrate the progress made in accelerator performance by the so called Livingston plot. It shows the maximum energy gain reached by installed accelerator systems since the beginning of the 20th century. This is visualized in figure 9. Here, in addition to the conventional, radiofrequency-cavity based accelerator systems (black dots), laser-plasma accelerators (red dots) and hybrid beam-driven accelerator experiments (black-red dots) are also plotted. It can be seen that with conventional accelerator technology, after the introduction of superconductor technology in the 1960s, the progress as regards maximum energy has significantly decelerated. In contrast, laser-plasma-accelerators, since the introduction of chirped pulse amplification in the second half of the 1980s, continue to advance at a much faster pace. With laser-plasma-accelerators, the size barrier is practically eliminated. With electric fields as high as approximately a TV/m, an accelerator of meter-scale size would be enough to break the TV-barrier. Improvement and refinement of schemes, in addition with constantly decreasing costs of laser technology, promise to continue this trend in the next years to come. To be fair, it shall be noted that the enhanced Livingston plot (figure 9) might be misleading to a certain extent, due to the limited comparability of the output of laser-plasma-accelerators when compared to conventional accelerators. Each type of accelerators has its advantages. For example, laser-plasma-accelerators are extremely versatile, and can, for example, produce electron beams with durations of a few fs, only, up to long beams with extremely broad bandwidth, whereas conventional systems can produce beams with extremely narrow energy spreads, and high repetition rates. Since the acceleration is much more rapid in plasma accelerators, electrons are much more suitable as particles to be accelerated when compared to protons and ions with their much higher inertia. In contrast, protons and ions can be accelerated easily in conventional ring accelerators, because their high masses makes them relatively insensitive to synchrotron radiation losses. It will be shown in the present ESA NPI project, that the high flexibility of laser-plasma-accelerators and the ability to produce broadband particle radiation are highly desirable characteristics as regards space radiation reproduction.

In the following, laser-plasma interaction will be discussed at some deeper physical level. Different laser systems which are or will be suitable for particle acceleration in the future will be presented. Then, the motion of electrons in intense laser pulse fields, ionization and absorption processes will be reviewed briefly. Finally, the different regimes of interaction, and the different radiation which can produced in these regimes, will be discussed.
Figure 9: Livingston plot enhanced by adding the performance of laser-plasma accelerators since the introduction of CPA. The maximum energy gain of laser- and beam-driven plasma-accelerators increases much faster than the energy gain in conventional accelerator systems, where the theoretical maximum accelerating electric fields have already been reached. It shall be noted that it is much harder to monochromatize the output of plasma-based accelerators due to the fundamentally different acceleration process when compared to conventional approaches.
4.1 Laser system technology fundamentals

This subsection sketches how high-power laser pulses are generated. Afterwards, these laser pulses can be focused and, provided proper choice of parameters, can then have focal intensities which induce nonlinear effects, up to being dominated by relativistic effects. The characteristic interaction of these focused laser pulses with matter is the driving force which results in release and acceleration of electrons as discussed in the upcoming sections.

4.1.1 Description of laser pulses

The propagation of electromagnetic waves are governed by Maxwell’s equations. In differential form and in vacuum these take the form

\[ \nabla E = 0 \]
\[ \nabla \times E = -\frac{\partial B}{\partial t} \]  
\[ \nabla B = 0 \]
\[ \nabla \times B = \mu_0 J + \epsilon_0 \mu_0 \frac{\partial E}{\partial t} \] (2)

This system of differential equations is coupled

\[ \text{rot rot } E = \text{grad(div } E) - \Delta E = -\text{rot } B = \epsilon_0 \mu_0 E \]
\[ \text{rot rot } B = \text{grad(div } B) - \Delta B = \epsilon_0 \mu_0 \text{rot } E = -\epsilon_0 \mu_0 B \] (3)

and describes the propagation of coupled \( E \) and \( B \) fields with the velocity of light \((1/c^2 = \epsilon_0 \mu_0)\) via the wave equation

\[ \left( \nabla^2 - \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \right) u = 0 \] (5)

A plane wave as a simple solution of this wave equation consists of time-dependent and phase coupled fields \((B_0 = E_0/c)\) which oscillate perpendicular to each other and to the propagation direction

\[ E(x, t) = E_0 \hat{y} \cos (kx - \omega t) \]
\[ B(x, t) = B_0 \hat{z} \cos (kx - \omega t) \] (6)

The energy density of this wave is composed by the energy densities of the electric field and the magnetic field

\[ w_{el} = \frac{1}{2} \epsilon_0 E^2 \]
\[ w_{mag} = \frac{1}{2\mu_0} B^2 \] (7)
resulting in a sum
\[ w_{\text{ges}} = \frac{1}{2} \epsilon_0 E^2 + \frac{1}{2 \mu_0} \frac{E^2}{c^2} = \epsilon_0 E^2 \] (8)

where \( B = E/c \) and \( (1/c^2 = \epsilon_0 \mu_0) \), or in the special case of a plane wave

\[ w_{\text{ges}} = \epsilon_0 E_0^2 \cos^2 (kz - \omega t) \] (9)

The intensity of a time-averaged electromagnetic wave, taking into account \( \langle \cos^2 kz - \omega t \rangle = 1/2 \), has the dimension of power divided by area

\[ \langle I \rangle = \frac{1}{2} \epsilon \epsilon_0 E^2 \] (10)

and is – in contrast to the electromagnetic fields – directly measurable. An extremely high energy density or intensity, respectively, is particularly attractive from a physical point of view and is the key to particle acceleration with lasers.

To achieve a high focused intensity, which is desirable for particle acceleration, the laser pulses must have high power \( P \). Various types of laser systems are known which can produce extremely high powers up to the PW regime. High pulse powers can be generated either by increasing the energy \( E \) of the pulse, or by decreasing the pulse duration \( \tau \), or both. In order to achieve extremely short pulse durations, the spectral bandwidth of the laser medium must be large. This is a direct consequence from Heisenberg’s uncertainty relation.

A laser pulse can be described analytically by modulating an infinite, monochromatic plane wave by a time-dependent envelope function \( E_A \), therefore by

\[ E(t) = E_A(t) \cos \omega_0 t + \phi(t) \]
\[ = 1/2 \left( \tilde{E}(t)e^{i\omega_0 t} + \text{c.c.} \right) \] (11)

where \( \tilde{E}(t) = E_A(t)e^{i\phi(t)} \).

Because due to Parseval’s theorem the spectrum of a pulse relates to the time domain via fourier transformation

\[ \tilde{E}(\omega) = \frac{1}{\sqrt{2\pi}} \int_0^\infty \tilde{E}(t)e^{i\omega t} dt + \text{c.c.} = \frac{1}{\sqrt{2\pi}} \int_0^\infty \tilde{E}(t)e^{i\omega t} dt \] (12)

the above pulse can be expressed by

\[ E(\omega) = E_A(\omega) \cos \omega t + \phi(\omega) \]
\[ = 1/2 \left( \tilde{E}(\omega)e^{i\omega t} + \text{c.c.} \right) \] (13)

where \( \tilde{E}(\omega) = E_A(\omega)e^{i\phi(\omega)} \).
Therefore, analogous to Heisenberg’s uncertainty relation, a larger spectral width $\Delta \omega$ implies a shorter pulse duration $\Delta \tau$ and vice versa

$$\Delta \tau \Delta \omega \geq \text{const.} \quad (14)$$

The quantity on the right hand side of the inequality is called pulse-bandwidth product, the value of the constant being dependant on the pulse shape. Because therefore an ultrashort pulse is intrinsically tied to a broad frequency spectrum, the dispersion, i.e. the phenomenon of the dependance of phase velocity of a wave of the frequency, must be considered and, if necessary, be compensated.

This is an extremely fundamental aspect of the generation of ultrashort laser pulses and reveals why Ti:Sa lasers with their exceptionally broad emission spectrum are so well-suited for high peak power pulse generation.

### 4.1.2 Single-shot amplifiers

Historically, Nd-doped glass lasers have been widely used for particle acceleration, and still are being used in low-repetition rate, highest energy and power systems. Prominent examples are the glass laser amplifiers used in the National Ignition Facility (NIF) in the US. Since the repetition rate for this system is low (few tens of shots per day), the inferior cooling characteristics of glass due to lower thermal conductivity when compared to Ti:Sa amplifiers is tolerable. At the same time, the advantages of glass amplifiers of being fabricable in large sizes at much lower costs can be harvested. The large sizes are needed in order to reduce the peak fluence of multi-kJ laser pulses at levels below the material thresholds.

### 4.1.3 Ti:sapphire amplifiers

Typical high-power laser systems used today for particle acceleration are solid state Ti:Sapphire lasers, which are based on a four-level system and emit laser light in the visible and near infrared spectrum. Due to the broadband energy levels in transition elements in general and in titanium-doted sapphire crystals in particular, not only the absorption frequency range which is responsible for optical pumping is very broadband, but also the emission range. Figure 4.1.3 shows the position and broadness of the absorption range in the green part of the spectrum and, articulately separated from this, the emission range in the red and near-infrared part of the spectrum.

In vacuum, the dispersion is given by the dispersion relation $\omega = ck$, $k$ being the wave number, the absolute value of the wave vector. Then, the phase velocity $\omega / k$ as well as group velocity $\partial \omega / \partial k$ of a wave packet with the frequency $\omega$ are always equal to $c$. In optical media, however, higher orders have to be taken into account, something that can be vividly seen from the splitting of light of several frequencies in a simple prism.
Figure 10: Ti:Sapphire absorption band in the green and broadband emission in the red part of the visible spectrum. Figure generated based on data from [22].

for example. Here, the phase velocity is different for different frequencies, resulting in different indices of refraction $\eta$. These indices of refraction are defined as the ratio of phase velocity in vacuum

$$v_{ph} = c = \omega(k)/k = 1/\sqrt{\epsilon_0 \mu_0}$$  \hspace{1cm} (15)$$
and in media

$$v'_{ph} = \omega(k)/k = 1/\sqrt{\epsilon \mu}$$  \hspace{1cm} (16)$$
that is by $\eta = c/v_{ph}$.

Using the wave number $k$, viz. the change of the spectral phase per unit length, the frequency dependency can be expressed as

$$k(\omega) = k_0 + \frac{\partial k}{\partial \omega}(\omega - \omega_0) + \frac{1}{2} \frac{\partial^2 k}{\partial \omega^2}(\omega - \omega_0)^2 + \frac{1}{6} \frac{\partial^3 k}{\partial \omega^3}(\omega - \omega_0)^3 + ...$$  \hspace{1cm} (17)$$

Whereas in zeroth order the phase velocity $k_0$, and in first order the reciprocal group velocity do not have an effect on the pulse shape, this is the case for higher orders. The second-order term is called Group Delay Dispersion (GDD) and influences the pulse shape due to its frequency dependence. With Quartz and a wavelength of 800 nm, this value amounts to $36 \text{fs}^2/\text{mm}$, for example. If the sign of this value is positive, as is the
case with most materials which are transparent in the visible region of light, this is called normal dispersion. If, in contrast, the term is negative, one speaks of anomalous dispersion.

With Quarz (SiO$_2$), the GDD at 1300 nm equals 0 – which is exploited in glass fibers, for example – and at even higher wavelengths negative. The TOD term (Third Order Dispersion) is frequency-dependent, too. It alters the spectral phase, which is why with ultrashort pulses for this order (and sometimes for even higher orders) measures have to be taken to compensate the dispersion.

### 4.1.4 Chirped Pulse Amplification Technique

Once an as short as possible and therefore spectrally broad laser pulse has been created in the oscillator, its energy can be increased in additional amplifier crystals. However, various effects such as thermal lenses, damage threshold of the amplifier material and other optical elements can hinder or limit the post-amplification. However, to a large extent, these effects can be dealt with very elegantly by applying the so-called Chirped Pulse Amplification (CPA) [23] scheme. Here, the laser pulse is getting stretched (for example by making use of above mentioned (positive) dispersion in optical media or by using grating systems) and accordingly, its power can be decreased by many orders of magnitude. Such a pulse, which still has exactly defined phase relationships after passing the stretcher, now can be amplified further without the danger of reaching the damage threshold. The amplified, temporally stretched pulse now can be re-compressed up to the theoretical Fourier bandwidth limit by using a complementary device, the compressor. Here, a system of gratings and/or prisms with opposed (negative) dispersion compensates for the stretching imposed before.

Such a broad bandwidth, ultrashort, highly energetic laser pulse can now in the last step be focused to enormous intensities making use of parabolic mirrors. The intensities producible by this scheme are many orders of magnitude higher than those normally occurring in nature and can therefore trigger extremely interesting physical effects. In the present work, these high intensities are being used to generate and accelerate pulsed particle beams for radiation effect testing.

In figure 11, the schematic principle of an appropriate laser system (here: a Ti:Sapphire laser system) is depicted and gives an overview on the essential steps which are necessary in order to produce such laser pulses.

### 4.1.5 OPCPA lasers

Next to Ti:sapphire and Nd:glass high-power-lasers, a third technique aiming at ultrashort, high-power laser systems is increasingly developed since recently. It is based on the Optical Parametric Amplification (OPA) technique, which can be traced back to the
early 1960s [24–28]. After CPA techniques became available in the late 1980s, it did not take long for the first realization of an OPA system combined with the CPA technique – the OPCPA laser [29]. Various optical birefringent crystals without inversion symmetry such as $\beta$-BaB$_2$O$_4$ ($\beta$-barium borate BBO) [30], KDB, LBO and much more [31] have been identified being suitable for OPA. In a well-cited publication of 1997, the prospect of using OPCPA lasers for amplification to multi-PW powers [32] has been anticipated. Today, actually most on-going or planned multi-PW laser systems aim at implementing high-gain OPCPA stages.

Good overviews on OPCPA techniques and laser systems can be found in [33–35]. Further, detailed information on OPCPA laser development can be found in [36] (TW-OPCPA laser development at Heinrich-Heine-University Düsseldorf, Germany) and [37–39] (OPCPA at MPQ Munich, Germany). In 2009, a sub-10-fs OPCPA laser system was for the first time used for quasimonenergetic laser-plasma electron acceleration [40].

The basic principle of OPA is strongly different from the Ti:sapphire amplification method. It is a second-order nonlinear process, based on the second-order nonlinear polarization (compare section 6, equation 33). Such a second-order polarization nonlinearity is also the basis for frequency doubling and sum and difference frequency generation. The highly simplified basic principle of OPA is visualized in figure 12. The pump beam

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**Figure 11:** Schematic summary of the processes which lead to the giant electromagnetic fields which are exploited by laser plasma acceleration.
and the signal beam propagate through the nonlinear crystal. The pump frequency (e.g., a frequency-doubled Nd:YAG laser at 532 nm) is higher than the signal frequency (IR). During the nonlinear interaction process, pump beam photons are converted into lower-energy signal and the same number of even lower-energy idler beam photons. Energy conservation is satisfied by the pump beam frequency being the sum of signal and idler frequency, $\omega_p = \omega_s + \omega_i$.

![Figure 12: Schematic visualization of Optical Parametric Amplification (OPA). The higher photon energy of the pump beam is converted into lower-energy signal as well as idler photons. This signal beam is amplified with very high gain.](image)

Noted advantages of OPA are the extremely high achievable single-pass gain ($\approx$ up to six orders of magnitude), as well as the gain bandwidth (supporting sub-10-fs pulses). Further, the thermal load on the amplifier medium is dramatically reduced, since the energy difference of pump and signal is not released into the amplifier (as in conventional solid state amplifier media), but is released as the idler wave, for which absorption in the crystal is negligible. The energy conversion is an instantaneous parametric effect. Also, the conversion efficiency can be very high, amounting up to about 20%. Finally, the parametric wave mixing process allows for extremely high tunability of the output signal. Next to the high gain bandwidth, wavelength are accessible which are otherwise not accessible directly with lasers. In the most simple, general case, wavelength tuning is produced simply by rotating the nonlinear crystal. These unique characteristics make OPCPA an extremely active candidate for future short-pulse, highest power laser systems. This said, the phase matching and other involved issues are highly complex as well as complicated, and are still subject to ongoing R&D.

### 4.1.6 Thin disc and fiber lasers

The attractiveness of using semiconductor diode lasers as pump sources for solid state lasers was recognized already in the 1960’s in the early days of laser technology. The high conversion efficiency both as regards the electrical to optical efficiency of semiconductor-based light emitting diodes (LED’s) as well as the high optical conversion efficiency from LED optical power to solid state laser power in addition to the size and cost reduction when using LED pumping was mentioned already in ref. [41]. Flashlamps
emit broadband radiation, and only a part of the emitted spectrum can be used to pump the laser, which limits the electrical-to-optical efficiency to about 1%. Semiconductor diode lasers, on the other hand, enable efficiencies of about 10%. In addition, the lifetime of diode-based pump lasers is about two orders of magnitude longer than those of flashlamps [42].

Fiber laser technology is one of the most flexible laser technologies, making it possible to operate in the cw as well as in the short-pulse mode, scalable powers and pulse energies ranging from the pJ to mJ level, and repetition rates up to hundreds of GHz. Since for laser-plasma-interaction in general, and laser-based particle acceleration in particular, high peak powers are necessary, we are interested in fiber laser technology which may lead to short pulses and ultrahigh powers. The very large spectral bandwidths of rare Earth element ions-doped glasses and other host materials make possible to access the few-femtosecond regime. Similar to solid state lasers such as Ti:Sapphire, this is the key towards ultrahigh peak pulse powers. Chirped Pulse Amplification (CPA) techniques are again necessary to reach highest energies without destroying the amplifier media.

Thin disc lasers are solid state lasers which use a thin disc-shaped laser crystal as the gain medium, the thickness of the disc (as thin as a few 100 µm) being smaller than the laser pulse diameter. A main advantage is that thermal load on the laser medium and associated detrimental effects such as thermal lensing can be minimized. The thin disc as the amplifier medium is often called an active mirror, because the one end of the thin disc (the cooled end on the opposite side of the pump and laser radiation) has a dielectric coating.

The advances of fiber lasers have been summarized in a Nature Photonics, Technology Focus issue in 2008 [43], and in [44]. For example, high repetition rate lasers based on a ytterbium-doped fiber chirped pulse amplification (CPA) system have been demonstrated to be able to produce compressed pulses as short as $\tau = 800$ fs with millijoule energies at a repetition rate of $\nu_{\text{rep}} \approx 100$ kHz, corresponding to an average power of more than 100 W. The performance of fiber lasers increased dramatically by about two orders of magnitude over the past decade [44], see figure 4.1.6. This is an unrivaled increase, even when compared to bulk solid state lasers.

One of the most important assets of fiber lasers are their extremely good wall-pug power conversion efficiencies into optical high-quality laser light. Since high brightness pump diodes in the 900-1000 nm wavelength became available at the end of the 1990s, ytterbium became the material of choice for high-power fiber lasers.

It is believed that fiber CPA systems have the potential for even shorter pulse durations, MHz rep rates and millijoule energies [45].

The pursuit of high power fiber lasers is further fueled by novel advances to combine multiple fiber lasers phase-coherently. For example, the International Coherent Amplification Network (ICAN) project aims at this, e.g. by combining 10000 laser pulses [46].
Figure 13: Increase of fiber laser power (cw) since 1996 (from [44]).
5 High-power laser systems in Europe

5.1 Introduction

Today, Europe has an especially strong position as regards high-power laser systems. In 2004, for example, two of the three groups which were reporting first generation of quasimonoenergetic electrons in Nature’s “Dream Beam” [47–50] issue were European. One was the group at the Laboratoire d’Optique Appliquée in France, using a laser system which delivered and energy of $E = 1$ Joule in a pulse duration of $\tau = 30$ fs, i.e. a pulse power of approximately $P \approx 30$ TW [50]. The other group used the Astra laser at Rutherford Appleton Laboratory in the UK, with an energy of $E = 0.5$ Joule in a pulse duration of $\tau = 40$ fs, corresponding to a power of $P \approx 12$ TW [48]. At that time, these lasers were the strongest available ones with pulse durations $< 50$ fs. These seminal results contributed greatly to the engagements of more and more groups in the field of laser-particle acceleration, and led to further mushrooming of high-power laser systems in Europe and in the world, since clearly various other groups recognized the potential and attractiveness of high-power laser systems for particle acceleration. It should be noted that a key goal in all these engagements was (and is) mainly the production of monoenergetic electrons (and, to a lesser extent, also of protons) with ultrashort pulse durations and high quality as regards divergence, emittance $\epsilon$, and energy spread.

To illustrate the dramatic increase in available high-power laser systems, it is helpful to concentrate on the development in Germany as an example. Here, in 2005, the only available high-power laser system was the JeTi laser at Friedrich-Schiller-University Jena, where at that time a peak power of about $P = 7$ TW was available, delivering an energy of about $E \approx 700$ mJ on target during $\tau \approx 80$ fs (at that time, the 2-TW ATLAS laser at the Max-Planck-Institute for Quantum Optics in Garching/Munich was undergoing a major upgrade). In contrast, in 2011, there were 5 laser systems with pulse powers $> 100$ TW available in Germany. The total available laser power of high-power short pulse ($\tau < 100$ fs) systems in Germany with powers $> 1$ TW increased from $< 10$ TW in 2005 to $\approx 1000$ TW in 2011, an increase of two orders of magnitude in approximately half a decade.

5.2 High-power laser laboratories in Europe

In Europe, several PW laser systems are already in use or planned. Among those are the Vulcan Petawatt laser, which is installed at the Central Laser Facility at the Rutherford Appleton Laboratory in Oxfordshire, England. Various laser beams are available at the Vulcan facility, among those a 700-fs, 700-J beam which can produce focus intensities of the order of $10^{21}$ W/cm$^2$. Next, there is PHELIX, a Petawatt High Energy Laser for Heavy Ion Experiments, installed at the GSI (Gesellschaft für Schwerionenforschung) at Darmstadt, Germany. There is another upcoming PW laser in Germany in Jena, which
is called Polaris, and which is based on the novel approach of using solid-state diodes for pumping of the amplifier medium instead of conventional flashlamp-powered pump lasers. The amplifying media in most PW-class lasers is Nd:glass, thus implying longer pulse durations (few ps down to a few 100 fs), and lower repetition rates (few shots per hour or minute) when compared to Ti:Sa lasers which are capable of producing pulses with durations down to about 10 fs due to their large spectral bandwidth, see equation 14.

European ultrahigh-power plans currently climax in the Extreme Light Infrastructure [51], a major project funded by the European Union with a final commitment exceeding 700 M€. This is the first large scale infrastructure based on the Eastern part of the EU. There will be four sites in total, the locations of the first three sites already having been determined. The site near Prague in the Czech Republic will focus on laser-plasma-particle acceleration, e.g. to produce ultrashort, fs-scale duration electron beams with energies of up to 10 GeV. The site in Szeged, Hungary, will focus on attosecond dynamics and ultrahigh intensity laser physics, while the third one in Magurele, Romania, will examine laser-based nuclear physics and generate an intense γ-ray burst by coupling a high-energy particle accelerator with a high-power laser. It is aimed at having these three sites operational by 2015. The location of the fourth site will be determined in 2012. This site aims at reaching laser pulse powers at the 200 PW level.

It shall be noted that a substantial part of ELI will be astrophysical research. For example, it is expected that ELI can contribute to the research on complex astrophysical environments, e.g. of pulsars, in the laboratory. It shall further be noted that ELI will promote an aggressive technology transfer. Both aspects should be highly interesting for European Space Agency as regards further engagement with laser-plasma-accelerator technology. The current ESA NPI activity could be a first link between ESA and high-power laser system research in Europe and in the member states.

Figure 14 gives an (non-exhaustive) overview of current multi-TW laser locations mapped on a map of Europe.

5.3 Laser collaborations

A major part of the highest-power laser system facilities in Europe are cooperating within the Laserlab Europe, an integrated initiative of European laser infrastructures, funded within the 7th EU framework [3]. It currently involves 26 laser research infrastructures from 16 EU member states. Figure 15 shows a mapped overview of laser facilities which are participating in Laserlab Europe. The facilities mapped in figure 14 and in figure 15 overlap largely, and complement each other. Laserlab Europe fosters various Joint Research Activities (JRA), among those HAPPIE (High Average and Peak Power lasers for Interaction Experiments) and LAPTECH (LAser Plasma acceleration TECHniques) possibly being the most important ones in the context of the present ITI activity. HAPPIE aims at developing novel laser technology which could be crucial to bring down the costs and size connected to laser-particle-accelerators suitable for space
Figure 14: (Non-exhausting) mapped overview of highest-power laser facilities in Europe.
radiation reproduction. One focus of the R&D work on Diode Pumped Solid State Laser (DPSSL) technology, for example, is the development of thin disc and fiber laser technology. Although commercial thin disc and fiber lasers are not yet available which have sufficiently large peak powers and focused intensities to be able to be of great interest for space radiation reproduction, there are experimental prototypes based on these techniques which reach the multi-mJ level and pulse durations $< 100$ fs, and are therefore highly suitable for the needs identified in the present ESA NPI project. So far, a general trend in high-power laser technology was that high-end experimental prototypes were transferred into the commercial world as comparably low-cost, reliable products in typically a few years at maximum.

![Figure 15: European laser facilities participating within the Laserlab Europe framework (figure from Laserlab Europe website [3]).](image)

### 5.4 Pulsed lasers with kHz and TW-level power

Next to the highest-power laser systems with repetition rates of 10 Hz or lower, a second class of Ti:Sa laser systems has been developed and is commercially widely available which is capable of kHz repetition rates. These lasers are in turn currently limited to
pulse energies up to 10 mJ. However, since these pulse energies can be delivered on target within pulse durations < 30 fs, and can optionally be further reduced to sub-8-fs making use of bandwidth increase via self-phase-modulation in rare-gas filled hollow fibers. Such laser systems are today often pumped with DPSS technology, with all the advantages for operationability such as reduced maintenance and increased power efficiency. Complete systems consisting of oscillator, amplifier, chiller and pump laser are currently being sold at the 300 k€ level. There are much more companies offering such systems when compared to the number of commercial manufacturers which are selling the highest-power multi-TW systems which have been described in the previous sections. These laser systems are mushrooming all over Europe and the world and are also usable for particle acceleration.

Figure 16: Femtopower Ti:Sa laser system with kHz repetition rate as installed at Heinrich-Heine-University Düsseldorf, Germany (figure adapted from T. Krüger, F. Gaussmann, private communication).
5.5 Commercial high-power laser systems

Concurrent with the demand from academia for high-power laser systems, industry recognized the developing economic potential of high-power laser systems for particle acceleration as well as as light sources. The pull from research was met by a corresponding industry push. As a result, today high-power laser systems are commercially available as turnkey systems from industry. In contrast, until a few years ago, all high-power laser systems worldwide were more or less internal developments mostly at research centers. Commercial highest power systems with pulse powers $> 100 \text{ TW}$ are currently available from two companies exclusively, which are both from France. The one is Thales Laser [52], with its world headquarter located near Paris, and the other is Amplitude Technologies [53], operated from close to Evry near Paris, together with its sister company Amplitude Systemes [54], located near Bordeaux.

Amplitude Technologies is offering, for example, Terawatt Systems with powers up to $P \approx 250 \text{ TW}$, with a compressed energy output of $E > 3 \text{ J}$ and a pulse duration of $\tau < 25 \text{ fs}$ at a repetition rate of 10 Hz. They fit in single laboratories with small footprints such as $\approx 30 \text{ m}^2$. Since approximately 2009, about four of these systems have been purchased and installed in Germany alone.

Very similar systems are offered by Thales Laser. In addition, Thales also offers laser systems with powers $P > 250 \text{ TW}$ up to 1.3 PW, at lower repetition rates of 0.1 Hz to 5 Hz. Still at pulse durations of $\tau < 30 \text{ fs}$, a peak power of 1.3 PW is reached with pulse energies of $E \approx 40 \text{ J}$. Such an 1.3 PW system is currently being installed for the BELLA project at Lawrence Berkeley National Laboratory, USA.

Next to these Ti:Sapphire lasers, thin-disk and fiber laser with repetition rates of 1 to 100 kHz are rapidly increasing in available pulse power [?]. Such systems are especially interesting and are believed to play a large role in the future also for radiation generation. They stand out by extremely good maintainability and wall-plug efficiency, along with other advantages such as the total cost of acquisition. A large and quickly increasing number of companies is active in this field.
6 Light wave interaction with matter

6.1 Interaction with single charged particles

In an electromagnetic field, the force on a charged particle is given by

\[ \frac{dp}{dt} = \frac{d(\gamma m_0 v)}{dt} = Ze(E + v \times B) \]  

where \( p \) and \( v \) are momentum and velocity, respectively, \( \gamma m_0 \) is the relativistic mass and \( Z \) the charge number of the particle.

As seen above, the magnetic field is much smaller than the electric field due to the \( 1/c \) factor, which means that for relatively slow (non-relativistic) velocities \( v/c = \beta \ll 1 \) the magnetic field term can be neglected. The oscillation momentum of an electron with \( Ze = e \) can be determined to

\[ p = -eE_L/\omega \]  

by simple integration in the time-dependent electric field. If this momentum is growing close to \( m_0c \approx 2.731 \times 10^{-22} \text{ kg m/s} \), the relativistic mass of the electron is getting significantly heavier than the rest mass, the motion of the electron in the laser field is getting anharmonic and, due to the then non-neglectable \( B \)-field contribution, a particle drift into the forward direction results. One speaks of relativistic oscillation velocities and introduces the so called dimensionless light amplitude

\[ a_0 = \frac{eE}{m_0\omega c} \]  

The parameter \( a_0 \) therefore separates the non-relativistic (\( a_0 \ll 1 \)) from the relativistic regime (\( a_0 \gg 1 \)). Since this decisive threshold is only dependent on the particle’s mass and laser frequency and intensity, the electromagnetic field amplitudes and the intensity can be written as

\[ E_0 = a_0 \frac{2\pi m_e c^2}{e\lambda} = \frac{a_0}{\lambda[\mu m]} \times 3.2 \times 10^{12} \text{ V/m} \]  
\[ B_0 = \frac{E_0}{c} = a_0 \frac{2\pi m_e c}{e\lambda} = \frac{a_0}{\lambda[\mu m]} \times 1.07 \times 10^4 \text{ T} \]  
\[ I = a_0^2 e_0 c \left( \frac{\pi m_e c^2}{e\lambda} \right)^2 = \frac{a_0^2}{\lambda^2[\mu m^2]} \times 1.37 \times 10^{18} \text{ W/cm}^2 \]

These equations are valid for electrons and a laser wavelength \( \lambda \) given in micrometers. For example, in case of a central laser wavelength of 800 nm, relativistic effects occur from a focused intensity of \( I \approx 2.14 \times 10^{18} \text{ W/cm}^2 \) on.

This intensity threshold is exceeded for nearly all experiments reported in this work. The movement of a single particle in an electromagnetic field is given by equation (18)
Figure 17: Single particle trajectories of free electrons in vacuum when interacting with laser pulses. a) Unfocused, non-relativistic laser pulse with $I = 2 \times 10^{16}$ W/cm$^2$ and $\tau = 9$ fs. b) Laser pulse as in a), but focused on a diameter of 3.5 $\mu$m (FWHM). c) Unfocused, relativistic laser pulse with $I = 5 \times 10^{19}$ W/cm$^2$ and $\tau = 80$ fs. d) As in c), but focused on a diameter of 3.2 $\mu$m (FWHM).
and can be calculated by integrating numerically. Figure 17 shows results of single particle tracking conducted using ELRAD [55] for several laser pulse intensities and laser focus sizes in vacuum. Figure 17 a) and b) are computed for the case of the interaction of a non-relativistic laser pulse with a pulse duration of $\tau = 9 \text{ fs}$, whereas c) and d) depict the trajectory of an electron experiencing a relativistic laser pulse with $\tau = 80 \text{ fs}$. The laser pulse in a) is not focused and has a homogeneous intensity of $I = 2 \times 10^{16} \text{ W/cm}^2$ according to $a_0 \approx 0.1$ on the complete lateral width. The electron oscillates during the few optical cycles of the 9 fs pulse only by few nanometers, in the end does not gain any kinetic energy and is located at the same lateral position as before the pulse. In contrast, in b) the laser pulse is focused on a diameter of 3.5 $\mu\text{m}$ (FWHM) and reaches the maximum intensity only in the focus. These laser pulse parameters correspond to the experiment using the D"usseldorf Femtopower system as described in section 3. In the simulation, the electron is located 1 $\mu\text{m}$ away from the axis and experiences a weaker restoring force in the second half of the laser cycle when compared with the first half due to the intensity gradient. All in all, a light drift away from the axis to the outside can be detected.

At non-relativistic intensities, which means when the Lorentz force can be neglected, the electron velocity in the alternating electric field is given by equation (18) and can be calculated by

$$m_e \frac{dv}{dt} = eE_0 \cos \omega t \quad (23)$$

which amounts to $v(t) = eE_0 \sin \omega t / (m_e \omega)$. Therefore, the kinetic energy is

$$\frac{m_e v^2}{2} = \frac{e^2 E_0^2}{2m_e \omega^2} \sin^2 \omega t \quad (24)$$

One defines the average energy of an oscillation movement as depicted in figure 17 a) as

$$U_p = \frac{m_e v^2}{2} = \frac{e^2 E_0^2}{4m_e \omega^2} \quad (25)$$

which via its gradient $\nabla U_p$ exerts a force on the particles. With $I = \epsilon_0 c E_0^2 / 2$ and hence $U_p = (e^2 \lambda^2 I)(8\pi^2 m_e \epsilon_0 c^3)$ one can give an easily applicable approximation for the ponderomotive potential as

$$U_p[\text{eV}] = 9.3 \times 10^{-14} \times I[\text{W/cm}^2] \times \lambda^2[\mu\text{m}] \quad (26)$$

For the example in figure a), the ponderomotive potential amounts about 1.2 keV.

The pulse in c) has an intensity of $I = 5 \times 10^{19} \text{ W/cm}^2$ according to an $a_0 \approx 4.8 > 1$ and therefore oscillates laterally by several 100 nm in the field of the 80 fs pulse, but as in a) returns to the laser axis after the pulse. In contrast, the pulse in d) is focused.
to a spot of 3.2 µm (FWHM). The laser pulse intensity profile is indicated by the gray curve. At the beginning of the interaction, the electron is 1 µm away from the laser axis, and oscillates further outside in the first half of the oscillation period of the laser pulse. The electron experiences relativistic mass increase and at the same time, owing to the magnetic field, a drift movement into the forward direction. Now the restoring force in the second half of the oscillation period is weaker due to the laser field which is decreasing to the outside and so on. As a result, this yields the plotted zigzag trajectory and the electron is effectively accelerated.

While the drift movement of charged particles in inhomogeneous electromagnetic pulses into the direction of the decreasing electric field was observed and interpreted already in 1957 [56], the name ponderomotive potential came into use not until considerably later [57, 58]. At higher, relativistic intensities such as in d) the ponderomotive potential (25), using $a_0$ and the relativistic $\gamma$ factor $\bar{\gamma} = \sqrt{1 + (p/mc)^2} \approx \sqrt{1 + a_0^2}/2$ averaged over one oscillation, takes the form [59–61]

$$U_p = -\frac{m_e c^2}{4\bar{\gamma} a_0^2}$$ (27)

Again, an engineering formula can be written as

$$U_p [\text{eV}] = -\frac{1.25 \times 10^5 \times a_0^2}{\sqrt{1 + a_0^2}/2}$$ (28)

hence, the ponderomotive potential in case d) (and the according experiments in section 5) can be estimated to about 0.85 MeV.

**Lawson-Woodward-Theorem**

As just seen, normally a plane electromagnetic wave cannot transfer substantial amounts of energy to a free electron. However, this is only valid under the following conditions:

1. the laser wave propagates in vacuum
2. the interaction is not limited and no boundaries are present
3. the electron is highly relativistic along the acceleration path
4. there are no additional static electric nor magnetic fields present
5. the $-v \times B$ force can be neglected.

This theorem is known as Lawson-Woodward theorem [62, 63]. In the following sections, this theorem is sometimes multiply infringed and substantial electron acceleration can set in.
In order to be accelerated in an electromagnetic field, a particle must be charged. Therefore, free charge carriers have to be released before. In case of the interaction of an electromagnetic wave with charge neutral matter the necessary process is ionization. The ionization is mostly caused by the electric field and can occur in several variants, dependent on the field strength. These are described later in this section.

6.2 Non-linear effects

At comparably low intensities before ionization starts, an electromagnetic wave with the electric field \( E \) does already influence the inner atomic fields and particles, resulting in the induction of a dipole moment and thus, a polarization

\[
P = \varepsilon_0 \chi_e E
\]  

(29)

where the so called susceptibility \( \chi_e = \varepsilon_r - 1 \) is a measure for the polarizability. This leads to a dielectric shift

\[
D = \varepsilon_0 E + P = \varepsilon_0 (\chi_e + 1) E = \varepsilon_0 \varepsilon_r E
\]  

(30)

In vacuum with \( \varepsilon_r = 1 \), the susceptibility is \( \chi_e = 0 \). Whereas in the regime of linear optics

\[
P = \varepsilon_0 \chi_e^{(1)} E
\]  

(31)

at higher field strengths and intensities, higher-order susceptibilities have to be taken into account

\[
P = \varepsilon_0 \chi_e^{(1)} E + \chi_e^{(2)} E^2 + \chi_e^{(3)} E^3 + \cdots
\]

(32)

which leads to nonlinear effects. Whereas the second-order susceptibility \( \chi_e^{(2)} \) vanishes for reasons of symmetry for most materials such as gasses, liquids, amorphous solids like glass and many crystals and is only effective in centrosymmetric crystals (such as KDP or BBO, which are used for second harmonic generation (SHG) for frequency doubling of laser light, for example), the third-order contribution \( \chi_e^{(3)} \) occurs in each material, in principle.

The non-linearity can be quantified by describing the electron bonding in the model of the (here, for simplicity, one-dimensional) anharmonic, damped oscillator in an electric field with Fourier components at the frequencies \( \pm \omega_1 \) and \( \pm \omega_2 \). This leads to the equation of motion

\[
\frac{d^2 x}{dt^2} + \Gamma \frac{dx}{dt} \omega_0^2 x + ax^2 = \frac{e}{m} \left[ E_1 (e^{-i\omega_1 t} + e^{i\omega_1 t}) + E_2 (e^{-i\omega_2 t} + e^{i\omega_2 t}) \right]
\]

(33)
Assuming the anharmonic term \( ax^2 \) in equation (33) is small at the beginning, one can use ansatz

\[
x = x^{(1)} + x^{(2)} + x^{(3)} + \text{c.c.}
\]  

and can get in the linear case

\[
x^{(1)} = x^{(1)}(\omega_1) + x^{(1)}(\omega_2) + \text{c.c.}
\]

solution

\[
x^{(1)}(\omega_1) = \frac{e}{m} \frac{E_1 e^{-i\omega_1 t}}{-\omega_1^2 + \omega_2^2 - i\omega_1 \Gamma}, \quad x^{(1)}(\omega_2) = \frac{e}{m} \frac{E_2 e^{-i\omega_2 t}}{-\omega_2^2 + \omega_1^2 - i\omega_1 \Gamma}
\]

Approximating the perturbation \( ax^2 \) by \( ax^{(1)}^2 \) yields terms at the sum and difference frequencies \( \omega_1 \pm \omega_2 \) of both waves the second harmonics \( 2\omega_1 \) and \( 2\omega_2 \) as well as at 0:

\[
x^{(2)} = x^{(2)}(\omega_1 + \omega_2) + x^{(2)}(\omega_1 - \omega_2) + x^{(2)}(2\omega_1) + x^{(2)}(2\omega_2) + x^2(0) + \text{c.c.}
\]

These terms are produced due to the quadratic character of the anharmonic term of the oscillator equation and can thus vividly explain the generation of sum and difference frequencies, the occurrence of higher harmonics as well as optical rectification [64]. In detail, the terms show the following dependencies [65,66]

\[
x^{(2)}(\omega_1 \pm \omega_2) \propto -2a \frac{e}{m} \frac{E_1 E_2}{\omega_0^6} \frac{1}{\omega_0} \frac{1}{\omega_0}
\]

\[
x^{(2)}(2\omega_{1,2}) \propto -a \frac{e}{m} \frac{E_{1,2}}{\omega_0^6} \frac{1}{\omega_0}
\]

\[
x^2(0) \propto -a \frac{e}{m} \frac{2}{\omega_0^6}
\]

For the electric polarization induced by an ensemble \( N \) it follows \( P = N e x \), which means that one can express the ratio of the contribution of the different order terms to the polarization from the summarized first-order terms \( x^{(1)} \) (see equation (36)) and the summarized second-order terms \( x^{(2)} \). In the non-resonant case \( \omega_0 \gg \omega_1 \) und \( \omega_0 \gg \omega_2 \) this is

\[
\left| \frac{P^{(2)}}{P^{(1)}} \right| \approx \left| \frac{a e E}{m \omega_0^6} \right|
\]

If the oscillation amplitude \( x \) of an electron orbital is so large that the linear force \( m\omega_0^2 x \) and the nonlinear force \( max^2 \) are getting similar, both have substantial effect on the position of the electron cloud. At the same time, they are then of the same order of magnitude of the inner atomic Coulomb force \( eE_{at} \) which is exerted on the electron by the nucleus, which means

\[
m\omega_0^2 x \approx max^2 \approx eE_{at}
\]
or by using the first and the second term

\[ \frac{m_e a^4}{a} \approx e E_{at} \]  

respectively. Inserting this leads to a ratio (41)

\[ \frac{P^{(2)}}{P^{(1)}} \approx \left| \frac{E}{E_{at}} \right| \]  

Analogously, one can also show that

\[ \frac{P^{(n+1)}}{P^{(n)}} \approx \left| \frac{E}{E_{at}} \right| \]  

enabling one to find useful estimations for the values of the susceptibilities of increasing order. In Bohr’s hydrogen model

\[ r_B = \frac{\epsilon_0 h^2}{(\pi m_e e^2)} = 5.29 \times 10^{-11} \text{ m}, \]

and with the elementary charge \( e = 1.602 \times 10^{-19} \text{ C} \) and the dielectric constant \( \epsilon_0 = 8.854 \times 10^{-12} \text{ As}/(\text{Vm}) \) one gets an electric field strength of

\[ E_{at} = \frac{e}{4\pi \epsilon_0 r_B^2} \approx 5.1 \times 10^{11} \text{ V/m} \]  

Linear optics tells us that \( \chi_e^{(1)} = \epsilon_r - 1 \) is of the order \( \sim 1 \). This yields for the first nonlinear order

\[ \chi^{(2)} = \frac{\chi^{(1)}}{E_{at}} \approx 2 \times 10^{-12} \text{ m/V} \]  

and for the next higher one

\[ \chi^{(3)} = \frac{\chi^{(1)}}{E_{at}^2} \approx 4 \times 10^{-24} \text{ m/V} \]  

and so on.

At the latest when the electric laser field grows larger than \( E_L = 10^{11} \text{ V/m} \), according to a laser intensity of about \( 10^{15} \text{ W/cm}^2 \), the series (33) does not converge anymore and thus is not applicable anymore. Here, the laser field is already big enough to induce field ionization.

### 6.3 Tunnel ionization and field ionization

At intensities higher than at multi-photon-ionization (MPI) or above-threshold ionization (ATI) levels, tunnel ionization is then the dominating ionization process. The transition can be described by the dimensionless, so called Keldysh parameter [67–69]

\[ \gamma_{\text{Keldysh}} = \omega_L \sqrt{\frac{2E_{\text{bind}}}{I_L}} \]  

\[ \frac{1}{4\pi} \]
in dependence of laser intensity $I_L$ and frequency $\omega_L$. With lower wavelengths and intensities $\gamma_{\text{Keldysh}} \gg 1$ and MPI/ATI dominates, whereas with increasing field strengths and wavelengths ($\gamma_{\text{Keldysh}} \ll 1$) tunnel ionization occurs predominantly. Here, the ionization rates can be determined by the so called Ammosov-Delone-Krainov theory (ADK) [70, 71].

Because the Ti:Sapphire laser systems frequently used for ionization and generation of plasma waves (which holds for the experiments in the present work, too) operate in the visible and near infrared range, and because often noble gases are used as media, the set of these wavelengths and the high ionization potentials of noble gases is particularly important. The field strengths of such lasers, where with noble gases the Keldysh parameter amounts to $\gamma_{\text{Keldysh}} \approx 1$, however, is so strong already, that field ionization (or barrier suppression ionization (BSI)) [72–75] sets in. Here, the potential barrier is so much decreased by the external field that quantum mechanical tunneling is not necessary anymore and the electron can leave the nucleus classically. Therefore, in case of the interaction of Ti:Sapphire laser pulses the MPI directly shifts to BSI.

As indicated in section 4.1.1 and following equation (10), an electric bonding energy of $E \approx 5.1 \times 10^{11}$ V/m such as in hydrogen can be overcome by a laser intensity of

$$I_{\text{ionization}} = \frac{1}{2} c \varepsilon_0 E^2 \approx 3.45 \times 10^{16} \text{ W/cm}^2$$

(50)

Figure 20 depicts the intensities and the according electric and magnetic fields in the laser focus over abroad range, and additionally and indicates with a couple of examples at which intensities additional physical effects occur. The red ellipse indicates at what intensities typical laser-plasma acceleration experiments take place.
Figure 19: First ionization energies of chemical elements in eV.
Figure 20: Visualization of laser pulse intensity and the corresponding light amplitude $a_0$ and the maximum electric and magnetic fields over a broad range. The shift from the non-relativistic to the relativistic regime occurs at $a_0 \approx 1$. 

Light wave interaction with matter

6.3 Tunnel ionization and field ionization
7 Laser-overdense interaction

The interaction of focused high-power laser pulses with overdense material (such as metal foils) can lead to efficient acceleration of electrons, protons, ions and other elementary particles. The interaction process can be broken down into three distinct components:

Interaction before the solid density: preplasma, laser absorption and reflection and electron acceleration

Propagation of electrons and ions into the solid

Emission of electrons and ions out of the solid

Matter is denoted "overdense" with respect to the driving laser pulse, if the electron density in the matter is so high that the charge displacement due to electron movement is so effective that the propagating electric field will be reflected. This density is called the critical density

\[ n_c = \frac{\omega_L^2 \varepsilon_0 \gamma m_e}{e^2} \]  

For Ti:sapphire laser systems, which have a central laser wavelength of \(\lambda_L \approx 800\) nm, corresponding to an angular frequency of \(\omega_L \approx 2.4\times10^{15}\) s\(^{-1}\), the critical density amounts to \(n_c = 1.7\times10^{21}\) cm\(^{-3}\). Matter in the solid state, such as metals, typically has a typical electron density of the order of \(n_e \approx 10^{22}\) cm\(^{-3}\), corresponding to a plasma frequency of \(\omega_p \approx 5\times10^{15}\) s\(^{-1}\) or a plasma wavelength of \(\lambda_p \approx 300\) nm. In experimental scenarios, there is always a more or less pronounced prepulse preceding the main laser pulse, for example due to amplified spontaneous emission (ASE) processes in the laser medium. As the increasingly intense pre-pulse sooner or later has focus intensities on target which are high enough to ionize matter, pre-pulses do generate pre-plasma. Typically, the plasma electron density \(n_e\) in this pre-plasma is exponentially decreasing from the solid surface into the vacuum, the scale lengths of this pre-plasma typically being the range of \(0.2\lambda - 5\lambda\), where \(\lambda\) is the laser wavelength. Coupling of the laser pulse energy with the surface takes place predominantly in the range where the electron density is in the range of \(n_c/4 < n_e < n_c\). Here, substantial parts of the laser radiation energy can be absorbed, which means energy is irreversibly transferred from the laser to plasma electrons.

7.1 Absorption processes

Absorption of laser pulses means irreversible energy transfer to particles. Generally, these processes are highly complex phenomena and mostly a mix of various mechanisms, whose contributions to the overall absorption often depend highly sensitively on the present situation. Nevertheless, in the following a rough classification of some of the most important absorption processes shall be tried.
Basically, the absorption of ultrashort laser pulses by matter can be separated into two big areas: one the one hand side, absorption primarily based on collisional effects and rather collisionless absorption on the other hand. Decisive parameters are particularly the sharpness of the density gradient at the vacuum-matter border, as well as the intensity of the laser pulse, and the resulting plasma or electron temperature, respectively. The transport of electrons in general, and their energy and temperature, respectively, in particular, play a dominant role, since they move directly in the laser field due to their low mass, and therefore here happens the primary energy transfer from non-ponderable energy (photons) to ponderable particles (electrons). For the induced electron movement and trajectories, the three decisive parameters are the ski-depth $\delta_S$, the electron dynamics in the laser field as well as the mean free path $[76]$.

Collisional absorption processes are

Inverse bremsstrahlung. Whereas conventional bremsstrahlung is based on the emission of a photon as the result of a collision of an accelerated electron with an atom or an ion, respectively, the energy absorption of an electron oscillating in the laser field and simultaneously undergoing collisions with atoms or ions is called inverse bremsstrahlung. Due to the collision, the electron escapes the mere oscillation in the laser field, which alone would not lead to a net energy gain. Inverse bremsstrahlung is especially strong when dealing with longer pulses, lower intensities ($\lesssim 10^{13} \text{ W/cm}^2$) and shorter laser wavelengths $[77]$. At higher intensities and higher electron velocities, respectively, the collision cross section and therefore the inverse bremsstrahlung absorption decreases, whereas its contribution increases with higher densities. The laser radiation damping rate amounts to $\nu \approx n_e \nu_{ei}/n_c$ and depends on the electron-ion collision frequency $\nu_{ei}$, which in turn scales with the number of released electrons per atom $Z$ as well as with the Coulomb logarithm $\ln\Lambda(\nu)$ $[78]$.

Inverse bremsstrahlung is very sensitive on the density profile close to the critical density, which is the range where the collision probability is especially high. The higher density which can be seen by a laser beam of higher frequency due to its deeper penetration ability, is also responsible for the increase of inverse bremsstrahlung absorption at shorter laser wavelengths. Inverse bremsstrahlung is a well-known effect $[79]$, since it plays a dominating role in laser fusion due to the relatively low intensities and long pulses which are used here.

(Normal) skin effect. In the extreme case of a perfect vacuum-metal border without any preplasma (hence describable by a Delta function) the laser pulse is abruptly incident on overcritical density.

Then the electric field $E_0$ incident on the vacuum-solid border decays with increasing distance $\zeta$ into the solid as

$$E(z) = E_0 e^{-z/\delta_S}$$

(52)
The exponential decay constant $\delta_S$ is called effective skin depth and amounts to

$$
\delta_S = \begin{cases} 
\frac{c}{\omega_p} & \text{wenn } \nu_{ei} \ll \omega_0 \\
\frac{c}{\omega_p} \sqrt{\frac{\nu_{ei}}{\omega_0}} & \text{wenn } \nu_{ei} > \omega_0
\end{cases}
$$

in dependence on the different damping scenarios.

Using the reflectivity $R$, the absorption $a$ can be written as

$$
a = 1 - R = \begin{cases} 
\frac{2\nu_{ei}}{\omega_p} & \text{wenn } \nu_{ei} \ll \omega_0 \\
\frac{2\omega_0}{\omega_p} \sqrt{\frac{\nu_{ei}}{\omega_0}} & \text{wenn } \nu_{ei} > \omega_0
\end{cases}
$$

The collisional skin depth is connected to the inverse bremsstrahlung, since the electrons oscillating in the evanescent laser field are losing their energy dissipatively via collisions.

Absorption processes which are to a large extent collisionless are

**Resonance absorption** [80]. At higher laser intensities ($\gtrsim 10^{15} \text{ W/cm}^2$) and correspondingly higher plasma temperatures the occurrence of collisions decreases and collisional absorption processes are hardly effective. In contrast to s–polarized light, p–polarized light (under non-perpendicular incidence) always has an electric field component in the direction of the density gradient into the target [81]. This field drives longitudinal plasma oscillations with a frequency $\omega_0$ into the target. At the point of the critical density $n_c$, this frequency is equal to the plasma frequency $\omega_p$, so that a resonance occurs and the plasma wave amplitude grows very big. For an obliquely incident light wave, which for reasons of simplicity oscillates in the x-y-plane, the dispersion relation is

$$
\omega^2 = \omega_p^2 + (k_y^2 + k_z^2)c^2
$$

Since the electron density varies only in one direction ($z$), $k_y$ is a conserved quantity. This means that $k_y$ is dependant on the angle of incidence $\theta$ such that $k_y = (\omega/c) \sin \theta$. It follows that the electromagnetic wave is reflected already in front of the critical density [82,83] at a density

$$
n_{\leftrightarrow} = n_c \cos^2 \theta
$$

Nevertheless, part of the electric field can extend up to the critical density (skin depth) and can excite the resonance. Resonance absorption can cause a laser pulse energy transfer as high as ca. 60 %. The effectiveness of resonance absorption is primarily dependant on the incidence angle, since this quantity determines the
distance between reversal point of the laser pulse and the location of the singularity. The bigger the distance, the smaller the evanescent field at the location of the resonance, since then it is harder for the incident laser wave to reach this point. On the other hand, the effect vanishes not only for big incidence angles → 90 (because a big \( \perp \)-component is important), but also for 0 (because then there is no propagation in the target normal direction). Therefore it is qualitatively clear that resonance absorption takes its maximum at an angle in between. The exact angle is dependant on the scale length

\[ L^{-1} = |d/dx \ln n_e(x)|_{x=x_c} \]  

(57)

and can be calculated quantitatively by using the the so called Denisov function \( \phi(\tau) \). Denisov’s function is a function of the dimensionless parameter \( \tau = (kL)^{1/3}\sin \theta \), being dependant not only on the angle \( \theta \) but also dependant on \( k \) (see above) as well as on the scale length \( L \) [80]. The absorption is then given as [84]

\[ a = \frac{\phi(\tau)^2}{2} \]  

(58)

The plasma wave itself can be damped by various mechanisms, which can be introduced via a damping term and which then cancel out the singularity in the wave equation. The can happen via various processes:

- Electron-ion collisions.
- Landau damping.
- Wave-breaking (also see section 8). Classical (cold) wave-breaking occurs, if elements of the plasma electron reservoir (which can be regarded as fluid) are equally fast (or even faster) than the phase velocity of the plasma wave. Electrons having this maximum velocity can be trapped by the wave and can be accelerated to high energies. The relatively few electrons are taking a good part of the laser energy, which appears as a high-energetic, suprathermal part in the electron energy spectrum [85]. The evanescent electric field which drives the plasma oscillation at the position of the critical density can be expressed as [84]

\[ E_{\text{drive}} = cB_c \sin \theta = \frac{\phi(\tau)E_0}{2\pi \sqrt{L/\lambda_0}} \]  

(59)

where \( B_c \) is the magnetic field at the position of the critical density and \( E_0 \) the electric field amplitude in vacuum. Assuming a damping collision frequency which is small when compared to the plasma frequency, then the absorbed energy is determined by

\[ I_{\text{abs}} = \frac{\pi \epsilon_0 \omega_0 LE_{\text{drive}}^2}{2} \]  

(60)
Now, conclusions can be drawn on the oscillation velocity of electrons accelerated in such way in the plasma wave. Due to thermal convection of the electrons with velocity \( \nu_t e \), the oscillation velocity of the electrons in the plasma wave has a maximum of

\[
v_{osc, wave} \approx 1.2 \left( \frac{\omega_0 L}{\nu_t e} \right)^{2/3} \nu_{drive}
\]  

(61)

where \( \nu_{drive} = eE_{drive}/(m_e \omega_0) \) is the oscillation velocity of electrons in the evanescent laser field \( E_{drive} \) at the resonance point [86]. The electron velocity necessary for cold wave-breaking to occur is [86]

\[
\nu_{break} = \sqrt{2 \nu_{drive} \omega_0 L}
\]  

(62)

A comparison of the necessary velocities for wave-breaking \( \nu_{break} \) and the oscillation velocity \( \nu_{osc, wave} \) shows that (using the thermal velocity \( \nu_t e = \sqrt{T/m_e} \)) that at temperatures \( T \) of a few 100 eV and at scale lengths up to a few 10 \( \mu \)m the necessary wave-breaking velocity is reached always. Therefore here, wave-breaking is the primary process and dominates clearly over the energy transfer due to collisions and Landau damping.

The excitation of a longitudinal plasma wave by the laser field, the expansion and breaking of this wave and the effective acceleration of electrons in this wave is one of the central topics of the present work. Under special circumstances, especially high-energetic electron pulses can be generated. In the case of the experiments with the Düsseldorf Femtopower laser system resonance absorption is not the process which leads to the highest energies.

Sheath inverse bremsstrahlung, anomalous skin effect.

Whereas inverse bremsstrahlung decreases with increasing electron energy, because then the collisional cross sections decrease, there is another, similar process which effectiveness complementarily increases with increasing electron velocity. Then the collision frequency decreases and the mean free path increases. Sheath inverse bremsstrahlung means that the mean free path is larger than the skin depth \( \lambda_{mfp} > \delta_S \), whereas the displacement of electrons in the laser field is smaller than the skin depth \( \nu_t e/\omega < \delta_S \).

In contrast, if the mean free path as well as the displacement of electrons in the laser field are larger than the skin depth \( \nu_t e/\omega > \delta_S \), one calls this situation anomalous skin effect, because then an electron can travel more than the skin depth during a laser cycle without the necessity of collisions and can therefore transport energy of the laser field deeper than the skin depth into the plasma.
The contribution of both effects to the total absorption is to a large extent independent on the laser intensity and amounts to an order of magnitude of about 10% [76].

Both effects are complementary to the so called vacuum or Brunel heating (see next point) in such a way as they occur primarily when the light pressure (due to the ponderomotive potential) is lower than the plasma pressure (the expansion due to the thermal energy of the electrons) [87]. The light pressure \( P_L = 2I_L/c \) increases with increasing intensity, whereas the plasma pressure \( P_P = n_e k_B T_e \) increases as temperature and electron density increase. If the light pressure and hence, the oscillation velocity of the electrons is smaller than the plasma pressure, which in turn is dependant on the thermal velocity, the effects at the plasma-vacuum border play a bigger role than at higher light pressures and thus, higher laser field strengths.

**Vacuum heating** (Brunel heating [88]). Next to resonance absorption, vacuum heating is one of the absorption processes which can be especially effective and can contribute by more than 50% to the total absorption. Like resonance absorption, vacuum heating is angle-dependant, because here the oscillation of the electric field in the direction of the target normal is the driving effect and angle-dependant itself.

At very strong light pressure or high intensities and sharp plasma gradients, respectively, we have \( v_{osc}/\omega > L \). Therefore, with \( p \)-polarized pulses under oblique incidence, under these circumstances electrons experiencing the laser field can be drawn out of the plasma (or via field ionization out of metals) into the vacuum in one half cycle and then can be turned back and accelerated into the plasma in the next half cycle. The order of magnitude of the velocity of the electrons which are reaccelerated into the target is of the order of the oscillation velocity

\[
v = v_{osc} = \frac{eE_L}{m_0\omega}
\]  

(63)

Since due to the short skin depth, the evanescent laser field does not have any effect on these electrons which are accelerated into the target in the next half cycle anymore, and thus the electrons deposit their energy via collisions in the high-density plasma or in the solid, respectively, and thus are able to heat the plasma effectively.

The effect is similar to the situation as described in section 6 and as depicted in figure 17: In vacuum and at nonrelativistic intensities a free electron would not have gained energy after the laser pulse has passed, whereas this is the case as soon as the laser pulses can not extert force uniformly on the electron during all half cycles. In this context, the ponderomotive force of a focused (relativistic) laser
pulse on a free electron does show certain parallels to the acceleration of electrons in high-density plasmas and the subsequent Brunel heating of the plasma.

Figure 21 shows schematically, how a linearly p-polarized laser pulse which is incident at an angle of $\phi = 45^\circ$ is reflected in the pre-plasma at the solid-vacuum boundary.

Figure 21: A p-polarized laser pulse is reflected at the vacuum-solid interface in the pre-plasma generated by the laser prepulse at near the critical density $n_c$.

In general, at long [84], but also at very short scale lengths the contribution of resonance absorption decreases. One can understand this behaviour quantitatively by the following reasoning: In the case of an extremely steep plasma gradient the electric field at the point of the critical density has not vanished yet and can be assumed to be the electric laser field, which is why the displacement of the electron elongation is $x_{osc} = \epsilon E_L/m_e \omega^2$. However, as soon as this elongation grows larger than the scale length $L$ on the vacuum side, the electrons cannot oscillate back anymore and the resonance breaks down.

Quantitatively, vacuum heating is mostly dominant for so called reduced scale lengths $L/\lambda < 0.1$ [89].

Numerous papers on absorption effects of high-power laser pulses impinging on overdense material have been published. A recent review can be found in [90].
7.2 Simulation of plasma generation and expansion

It shall be noted again that the electron temperature as well as the emitted charge is furthermore dependent on the preplasma, and thus the laser prepulse. Preplasma formation can be estimated using hydrodynamic tools such as MULTI-FS. MULTI-FS [91] is a 1D version of the code MULTI [92], which offers the advantage to be able to describe absorption and plasma expansion not only on the ps-level but also on the fs-timescale.

In order to achieve this, essentially three changes have been implemented in MULTI-FS:

- assumption of thermal equilibrium between electrons and ions with different Equations Of State (EOS) for electrons and ions (instead of thermal equilibrium with the same EOS as in MULTI),
- modelling of electron collisions and the resulting absorption of laser light in the interaction zone, thermal conduction by electrons moving into the dense target, as well as energy transfer between electrons and ions,
- the propagation of the incident electromagnetic wave in a steep plasma gradient by solving Maxwell’s equation in a fine mesh.

The separate treatment of electrons and ions is essential, since during short-pulse interactions it is important that at first the electrons are heated. The relaxation time into the thermal equilibrium between electrons and ions is of the order of only a few to several 10 ps, which means that at first the electron temperature is many orders of magnitude higher than the ion temperature.

The energy of ultrashort femtosecond laser pulses primarily is being deposited in volumes with steep density gradients ranging from critical density as far as up to solid density (compare figure 21). As a basic principle, an electromagnetic wave can propagate through the medium if it is underdense. When an electromagnetic wave is incident on a plasma, under certain circumstances plasma electron oscillations are excited. Due to the much higher inertia of the ions one can assume that in the oscillating electromagnetic field, at first only the electrons move. Due to the quasistatic ion background there is a restoring force \( F = eE \). Making use of basic electrostatic laws and Poisson’s equation, the electric field \( E \) which is dependant on the displacement can be calculated via a simple integration, additionally making use of the electron density \( n_e \), which is also displacement-dependant. One yields a harmonic oscillator with the plasma frequency

\[
\omega_p = \sqrt{\frac{n_e e^2}{m_e \varepsilon_0}}
\]  \hspace{1cm} (64)

Now the oscillation of the electrons in the laser field leads to a situation where at a low enough oscillation frequency of the incident electromagnetic wave, the electron density is sufficient to make the electrons follow the perturbation fast enough in order...
Laser-overdense interaction

To shield the electromagnetic wave and thus stop it from propagating any further, the critical frequency $\omega_c$ is the result of the dispersion relation for electromagnetic waves in plasma

$$\omega_c^2 = c^2 k^2 + \omega_p^2$$

(65)

Therefore, as soon as the laser frequency $\omega_L$ gets smaller than the plasma frequency $\omega_p$, this would mean that $k$ would have to be imaginary, which physically manifests by that the wave cannot propagate any further. From the plasma’s perspective, this point is exactly where the plasma frequency equals the laser frequency, which in terms of density is at the critical density

$$n_c = \frac{\omega_L^2 \epsilon_0 \gamma m_e}{e^2}$$

(66)

In the hydrodynamic computer code, the range with the steep density gradient is divided into narrow slices, the characteristics of which are reflected via the index of refraction $\eta$ and the dielectric constant, respectively [93]

$$\epsilon = \eta^2 = 1 - \frac{\omega_p^2}{\omega_L (\omega_L - i\nu_{ei})}$$

(67)

This equation contains the laser and plasma frequencies $\omega_L$ and $\omega_p$ as well as the electron-ion collision frequency $\nu_{ei}$. This collisional frequency also spans many orders of magnitude, due to the extremely high density and temperature range in the interaction zone of laser and target. Therefore this range has to be accessible by the code. MULTI-FS uses a theoretical model which is well portable to a computer code. For reasons of simplicity, it is based on the two limiting cases of the collisional frequency in the cold solid on the one hand and on the collisional frequency in the hot, ideal plasma on the other. In between, the code interpolates cleverly [94].

Exemplary simulations with MULTI-FS have been carried through in order to determine the influence of prepulses of different intensities on the formation of a preplasma for an angle of incidence of the laser pulse of 45°, p-polarized light and a sin² laser pulse. The target material was aluminium, for which the EOS are especially well-known. In all simulations, a pulse duration of 25 fs was chosen, reflecting a typical pulse duration of commercial high-power laser systems which are today used for laser-plasma-acceleration. The code has been modified in such a way as to yield the electron and solid density at the vacuum-target interface at various points of time after incidence of the laser pulse.

Figure 22 depicts the temporal development of the electron (a) and solid (b) density for a laser pulse with a wavelength of 800 nm which is incident from the right and has an intensity of $10^{14} \text{W/cm}^2$ after $t = 15, 50$ and $100$ fs. Analogously, figure c) and d) show the plasma expansion for a prepulse of intensity $10^{15} \text{W/cm}^2$. The dashed line indicates the critical density.

From the simulation results it can be learned that the scale length amounts – even for a peak intensity of $10^{15} \text{W/cm}^2$ – to only a few nm within the time span of interest of
Figure 22: Hydrodynamic 1D simulations with MULTI-FS for a hypothetical lase prepulse (coming from the left, 25 fs duration, target material: Al). a) and b) show the temporal development of electron and material density for a peak intensity of $10^{14}$ W/cm$^2$. Analogously in c) and d) for $10^{15}$ W/cm$^2$. 
up to 100 fs. In fact, the scale length of the preplasma determined by the hydrodynamic simulations is even thinner than the skin depth of the laser pulse in the solid (aluminium), in other words the depth of penetration, after which the electric field of the laser pulse has decreased to the 1/e-fold [83, 95]. For the same laser pulse, this was confirmed via experiments [96, 97].

One the one hand this means that the underdense range in the plasma, which the incident laser pulse has to propagate through up to the critical density, is very thin, and in consequence this plasma plays an inferior role during the interaction of the pulse with the target. In contrast, a relativistic prepulse would result in a substantially thicker preplasma. The plasma expansion resulting from a prepulse with an intensity of $10^{18}$ W/cm$^2$ and the generation of a shock front propagating into the solid is shown in figure 23.

Particle-in-cell simulations can be used to simulate the outcome of the interaction of the main laser pulse with the pre-formed plasma. Ideally, they would use as input the results of hydrodynamic simulations as those described above in order to model the preplasma conditions at the moment of the incidence of the main pulse right. This said, it should be noted that PIC-simulations of laser-solid interactions mostly aim at predicting the acceleration of secondary particles such as protons or ions. In this context, the real solid-density plasma is often not modeled accurately, but in order to simulation time and computational load, often a reduced density (e.g., ten times overcritical) electron density is used.

Generally, in laser-plasma interaction one aims at an as low prepulse as possible (again, for example for certain proton and ion acceleration scenarios). However, sometimes a stronger pre-pulse can be advantageous, since it increases the preplasma which is present when the main pulse arrives. Therefore, sometimes a prepulse is even deliberately
generated.

7.3 Electron acceleration

7.3.1 Fundamentals

By various mechanisms, for example by those similar to as shown in section 6 and figure 17, electrons can now be accelerated in the pre-plasma in the forward and target normal direction. Because of the oscillating character of the laser wave, electrons are accelerated in a broad range of energies. Only few electrons are accelerated to the highest kinetic energies $E_{\text{kin}}$, and the stochastic nature of the acceleration process typically leads to electron energy spectra which can be approximated by an exponential decays function.

The electron number distribution typically follows a so-called exponential spectrum

$$N \propto \exp \left( \frac{E_{\text{kin}}}{k_B T} \right)$$

where $k_B$ is Boltzmann’s constant, $T$ is the electron temperature in Kelvin, and $T_{\text{eff}} = k_B T$ is the effective electron temperature in eV.

The interaction of focused, high-power laser pulses with solid target foils is a very typical scenario. In this context, "engineering" scaling laws have been developed which predict the electron temperature of electron beams generated during laser-solid interaction in dependence of the focused laser intensity. Work on this had already started in the era before Chirped Pulse Amplification [85, 98, 99]. After laser powers increased through CPA [23] techniques (see section 4.1.3) techniques, these scalings have been further developed towards relativistic laser intensities [100]. In general, the effective electron temperature can be described with a $T_{\text{eff}} \propto (I \lambda^2)^\zeta$ scaling, where the exponent $\zeta$ typically has values between $1/2$ and $1/3$ [89]. In the scaling of Wilks [101], it was found that in the intensity range from $I \lambda^2 \approx 1.3 \times 10^{18}$ W$\mu$m$^2$/cm$^2$ to $I \lambda^2 \approx 1.4 \times 10^{19}$ W$\mu$m$^2$/cm$^2$ the electron temperature scales as $\sqrt{T}$. The scaling is only dependent on the laser intensity and the laser wavelength and reads in its explicit form

$$T_{\text{eff, Wilks}} = \left( \frac{1 + I [\text{W/cm}^2] \lambda [\mu\text{m}]^2}{(1.37 \times 10^{18})} - 1 \right) m_0 c^2.$$  

Another scaling, which is based on experiments with sub-ps laser pulses and intensities up to the $10^{19}$ W/cm$^2$ level and which follows a $I^{1/3}$ scaling [102] was developed by Beg. It differs from Wilks’ scaling in that it predicts higher temperatures for intensities lower than $\approx 2.8 \times 10^{18}$ W/cm$^2$ and lower temperatures than Wilks for higher intensities. Explicitly, it is

$$T_{\text{eff, Beg}} = 0.1 (I_{17} \lambda^2)^{1/3} \text{MeV}$$

where $I_{17}$ gives the intensity in multiples of $10^{17}$ W/cm$^2$. A new analytical approach to the electron temperature was made recently by Kluge et al. [103] which fits the
experimental and numerical data in the range of $a_0 \approx 5$ as good as the Beg-scalings and for $a_0 < 1$ fits to the ponderomotive scaling. Thus, by changing the intensity, and by changing the preplasma parameters, one can easily steer the temperatures of the generated electron beams. Figure 24 visualizes the electron temperature in dependence of the laser focus intensity as predicted by the scalings of Wilks [101] and Beg [102].

Figure 24: Expected electron temperature according to the scalings of Wilks [101] and Beg [102]. Beg’s scaling is better suited for low intensity, Wilks’ scaling is more appropriate for higher, relativistic intensities.

For similar physical reasons which lead to broadband, exponential energy spectra of the generated electron beams, the electron beams are emitted in a broad cone into the forward direction, with an opening angle of approximately $30^\circ$. The divergence of the electron beam is also intensity-dependent, which has been published in a comparative publication which analyzed in 2008 [104] the divergence measured in various laser-solid interaction experiments. Most of these experiments were based on laser pulse durations between 350 and 500 fs [105–109] in the context of laser-based fast ignition, and were complemented by results from shots with 5 ps duration pulses [104]. In these experiments, the intensity ranged from $10^{18}$ W/cm$^2$ to $10^{21}$ W/cm$^2$. These data points were further complemented with measured results from experiments with 80-fs laser pulses [110]. The basic result is that there is no significant dependence of the observed electron temperature when the pulse duration is varied, but that there is a clear and almost linear increase of electron temperature as the laser intensity increases. These results are summarized in figure 25.

On a theoretical basis, the divergence of laser-driven relativistic electron beams was explored in [111]. The divergence when using cone targets, which are highly relevant for fast ignition inertial confinement laser fusion schemes, was examined in [112].

Summing up, the divergence as well as the electron temperature can be steered by varying the laser intensity. Varying the laser intensity, in turn, can be by changing...
the laser pulse energy, or, even simpler, by moving the target out of the laser focus. Depending on the strength of focusing (i.e., the F-number of the optical element used to focus the laser beam), shifting the target foil out of focus by tens of microns can already be enough to decrease the effective intensity on target substantially. The important parameter here is the Rayleigh length $z_R = \pi \omega_0^2 / \lambda$, where $\omega_0$ is the beam waist. The beam spot size of a Gaussian beam at a distance $z$ away from focus is

$$\omega(z) = \omega_0 \sqrt{1 + \left(\frac{z}{z_R}\right)^2}$$  \hspace{1cm} (71)

and the peak intensity when varying the distance $z$ amounts to

$$I(z) = \frac{2P_0}{\pi \omega^2(z)}$$  \hspace{1cm} (72)

where $P_0$ is the peak power of the laser pulse. It shall be noted that if the intensity is varied by moving the target out of focus by a distance $\Delta z$ (without increasing the laser energy or power $P_0$), not only the intensity decreases according to equation 72, but also the spot size increases according to equation 71. At the moment, no extensive data exists which predicts how the total electron charge which is accelerated per shot scales with the spot size. However, it is reasonable to assume that the total charge scales linearly with the laser spot size area on target. Experimental data therefore needs to

Figure 25: Measured divergence of electron beams produced in various laser-solid-interaction experiments. As figure a) shows, there is approximately a linear dependence of the laser intensity, while there is no clear connection to the pulse duration as shown in b). Data from: Stephens [105], Santos [106], Lancaster [107], Kodama A [108], Kodama B [109], Green A [104], Green B [104], Hidding et al.
be assembled to explore the relation between spot size and emitted electron charge at constant intensity. Typically, electron charges of the order of tens of nC are emitted per shot.

7.3.2 Electron propagation through the solid target

Up to $\approx 40\%$ of the laser energy can be absorbed by the plasma generated at the vacuum-solid interface. Most of this energy goes into hot electrons which are pushed forward into the solid almost collisionless. This hot electron current is being (quasi-)compensated by return currents, which are much colder and therefore determined by collisions.

The necessity of such return currents can be accounted for in different ways, for example

- via the Alfvén limit.
- by an energy balance.
- by self-generation of electric and magnetic fields.

These three explanations will be briefly discussed in the following.

Alfvén limit

It is well-known from the field of astrophysics [113], that jets of charge carriers are subject to certain restrictions as concerns their maximum current. Generally speaking, the so-called Alfvén limit is that amount of current, at which point any hypothetically added charged particle would experience a net backward drift due to its trajectories under the influence of the self-generated magnetic field of the whole current. The particle would be bent back to its relative origin. The value of this basic Alfvén limit amounts to

$$I_A \approx \frac{4\pi\varepsilon_0 mc^3}{e}\beta\gamma \approx 17\beta\gamma \text{kA} \quad (73)$$

As already pointed out by Alfvén, this is strictly valid only for a special case, nevertheless giving the order of magnitude of this current restrictions correctly for many other scenarios [113, 114]. Indeed in more realistic and complicated scenarios – for example when the spatiotemporally structure of the pulsed current is important – one has to deal with variety of highly complex of (sometimes competing) effects, which is why there is a lively debate on all kinds of aspects of Alfvén limitation of currents up to today [114–118]. This debate recently gained even more in importance due to the upcoming principle of fast ignition. However those currents needed for fast ignition, for instance, are much higher than the estimated Alfvén limit [114], and can indeed be reached with laser-plasma interactions in practice.

The number of electrons generated in such an interaction with a Boltzmann-like energy distribution population can be estimated by the ratio of the laser energy transferred and
the average energy of the electrons. For example, from experiments carried through with a 80-fs, 800 mJ laser pulse [9], it can be estimated following [119], that as a lower limit a fraction of $\eta \approx 20\%$ of laser energy $E_L \approx 800 \text{ mJ}$ has been converted to electrons with a temperature of about $k_B T \approx 1.5 \text{ MeV}$.

The current which at first propagates into the target therefore does contain about

$$N_e \approx \frac{\eta E_L}{k_B T} = 6.7 \times 10^{11} \quad (74)$$

electrons. These are generated in a period of time of the order of magnitude of the laser pulse duration of about $\tau \approx 100 \text{ fs}$. The resulting average current therefore has a strength of about 1 MA. This amount is at least one order of magnitude higher than the Alfvén limit and can only be reached, because there are return currents which counteract and diminish the magnetic field created by the forward current.

**Energy balance**

Applying Ampère’s law known from magnetostatics

$$\oint \mathbf{B} \, ds = \mu_0 I_{\text{ges}} \quad (75)$$

one can estimate, following a reasoning from [120], under the estimate of a total current of $I_{\text{ges}} = 1 \text{ MA}$ propagating inside a cylinder with a diameter of the order of the laser pulse spot size $r_{\text{spot}} \approx 5 \mu\text{m}$, that the magnetic field at the cylinder surface is

$$B_{\text{cyl}} = \frac{\mu_0 I_{\text{ges}}}{2\pi r_{\text{spot}}} \approx 40 \text{ kT} \quad (76)$$

Using the energy density

$$u_B = \frac{B_{\text{cyl}}^2}{2\mu_0} = \frac{\mu_0 I(r)^2}{8\pi^2 r^2} \quad (77)$$

and the radius-dependent current

$$I(r) = \frac{I_{\text{ges}} \pi r^2}{r_{\text{spot}}^2} \quad (78)$$

and a propagation distance of electrons into an arbitrarily long target of $R_e \approx 100 \mu\text{m}$ [120] there is an integrated magnetic energy $U_B$ of

$$U_B = R_e 2\pi \int_0^{r_{\text{spot}}} \frac{\mu_0 I_{\text{ges}}^2}{8\pi^2 r_{\text{spot}}^2} \, r \, dr \quad (79)$$

$$= \frac{R_e \mu_0 I_{\text{ges}}^2}{4\pi r_{\text{spot}}^2} \int_0^{r_{\text{spot}}} \, r \, dr \quad (80)$$

$$= R_e \frac{\mu_0 I_{\text{ges}}^2}{8\pi} \quad (81)$$

$$\approx 5 \text{ J} \quad (82)$$
However, this is energetically impossible, since with the invested laser energy in underlying experiments amounted to „only“ 0.8 J. This paradox is being resolved, if one allows return currents which have an opposing effect on the magnetic field generation and therefore decrease its energy.

**Self-generation of electric and magnetic fields**

The laser pulse duration and the duration of the electron pulse driven by this laser pulse is much longer in a solid than the plasma period $\tau = \frac{2\pi}{\omega_p}$, due to the high density. Therefore the solid has to be close to charge neutrality, the charge compensation happening within the very small collision-free skin depth $c/\omega_p$ [121]. This could be confirmed by PIC simulations [122, 123]. The timescale for charge and current neutralization can be given as $\tau_n = \epsilon_0 \rho$, where $\rho$ is the specific resistance [124]. Even with a typical limit of the specific resistance in conductors of $\rho \approx 2 \mu \Omega m$, the value of $\tau_n$ amounts to only 17.7 as.

The (to a large extent) charge neutrality therefore requires quasi-compensating, comparably cold return currents, which due to the small skin depth spatially coincide with the hot, directly laser-generated forward electron beam. Therefore, locally, there has to be quasi-compensation of the current densities [120, 121, 125, 126]

$$j_{\text{hot}} + j_{\text{return}} \approx 0$$  \hspace{1cm} (83)

The counterpropagating, cold electron current is resistive and creates an electric field which decelerates the forward current of hot electrons (electric inhibition)

$$E = \eta j_{\text{return}} = -\eta j_{\text{hot}}$$  \hspace{1cm} (84)

However, a system of counterpropagating currents is unstable and can result in breaking-up and filamentation of the fast, relativistic electron pulse [122, 123, 127–129] due to the Weibel instability [130], which could also be observed experimentally during laser-plasma interaction [131].

Not only the hot current is connected with magnetic fields via

$$\frac{dB}{dt} = -\nabla \times E = \nabla \times \eta j_{\text{hot}}$$  \hspace{1cm} (85)

but also the cold currents, which means that not only the currents, but also the magnetic fields compensate each other.

However, since the current neutralization is ultrafast, as discussed above, but can never be instantaneous, the compensation of magnetic fields is neither perfect. Due to the huge currents a small fraction of non-compensated current is enough in order to create a huge magnetic field. As estimated by using equation (76), the quasistatic magnetic field can reach e few 10 kT in case of the experiments done with the JETI laser. Such a magnetic field is of the same order of magnitude as the oscillating $B$-fields
in the focused laser pulses. These $B$-fields as well as the quasistatic $B$-fields generated by the forward-propagating electrons, are substantially higher than all other man-made magnetic fields so far created in any laboratory. For example, at a laser intensity of $I_L \approx 5 \times 10^{19} \text{W/cm}^2$ as in the experiment, the maximum of oscillating magnetic field of the laser pulse amounts $B = 64 \text{ kT}$, which is higher, but of the same order as the magnetic field created by the laser-generated brutto current, which amounts to about 1 MA (see section 7.3.2).

If one considers the scenario of a break-up of this current into filaments each with a maximum current as given by the Alfvén limit, the magnetic field, using equations (73) and (76) with a filament radius of about $5 \mu\text{m}$, would amount to 'only' about 10 kT, but nevertheless the magnetic fields in turn can lead to a constriction of the filaments [122, 123, 127–129], which means that principally there is the possibility that their radii are temporarily smaller than the laser focus, again increasing the maximum according magnetic fields to higher values.

### 7.4 Accelerated electrons emitted into the vacuum

After being driven in the forward direction under the influence of the $v \times B$ force, due to Brunel heating etc., and after their complex, field-determined propagation inside of the target under the influence of return currents, a fraction of the electron beam can leave the target at its rear side. Mainly the slower electrons are responsible for creating electrostatic fields at the target foil rear side, which are in turn responsible for the secondary acceleration of protons and ions in the TNSA principle (Target Normal Sheath Acceleration) [132]. The faster part of the electron population is fundamental for the fast ignition scheme (in addition, there are considerations to use the protons produced via TNSA for ignition in the fast ignition scenario [133]).

Measurement of the electrons emitted in the experiments at the target rear side can be done using various diagnostics. Although in principle diagnostics as known from state-of-the-art accelerators can be used, such as beam profile monitors (BPMs) etc., in the special case of laser-plasma-accelerators the diagnostics should be capable to deal with an especially harsh radiation environment. Beam viewers and calibrated image plates are suitable for such an environment and are therefore among the most important diagnostics.

In order to quantize the electrons emitted into the forward direction as regards their energy spectrum a straightforward way is to use the energy-dependent deflection in magnetic fields generated by permanent magnets. The design and construction process can be performed by using state-of-the art tools such as the CST PARTICLE STUDIO code, which belongs to a code framework which is often used in accelerator physics.

Electron energies obtainable from standard laser-solid interaction with state-of-the-art Ti:Sapph lasers in the multi-TW regime have Boltzmann-like spectra with cutoff-energies not over a few tens of MeV. Therefore, it is neither necessary nor useful to use
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Neodymium-iron-boron (NdFeB) based permanent magnets with a magnetic remanence value of $B_r \approx 1.1$, which are among the strongest commercially available magnets (these are necessary to resolve GeV-scale electrons which can be produced via underdense interaction, see section 8). Instead, hard ferrite magnets based on barium ferrite and strontium ferrite (BaFe, SrFe) with a magnetic remanence of $B_r = 380 - 400 \text{ mT}$ can be used. In connection with a U-shaped yoke made of magnetizable steel to bundle the magnetic flux a to a large extent homogeneous magnetic field for deviation of the electrons can be generated.

For example, figure 26 shows the magnetic configuration of such a compact spectrometer based on a pair of hard ferrite magnets sized $60 \text{ mm} \times 20 \text{ mm} \times 15 \text{ mm}$ with a gap between both magnets of $10 \text{ mm}$, which generates a magnetic flux density of about $0.23 \text{ T}$. Figure 26 a) and b) show in detail the configuration of the spectrometer as well as simulation results of the generated $H$- and $B$-fields. The calculated values of $B$ have additionally been experimentally spot checked by means of a Hall sensor, showed excellent agreement with the simulations.

Figure 26: Simulation results for the magnet-based spectrometer for electrons up to 30 MeV. a) Structure of the $H$-field created by the twin magnet and the yoke. b) Same for the $B$-field. The effective field between the magnets amounts to about $0.23 \text{ T}$. c) Calculated trajectories of electrons with energies ranging from 0.5 to 30 MeV passing through this magnetic field. d) Energy calibration curve.
Due to this field between both magnets, electron which are incident on the laser axis are now deviated according to the Lorentz force and can then be detected energy-resolved. In 26 c), calculated trajectories for electrons passing the aperture are plotted, and in d) the corresponding calibration curve for the points of incidence of different electron energies in the detection plane is given.

In the detection plane (which is parallel to the laser axis in the depicted case), either electron-sensitive phosphorescent screens such as those described in [134] or electron-sensitive image plates such as those described in [135] (and references therein), or a combination of both, can be used. In case of a phosphor screen, the electron spectrum can be retrieved online by making use of a triggered camera observing the screen, whereas in order to read out absolutely calibrated image in a special scanner device it is necessary to break the vacuum.

The front of such an electron spectrometer has to be radiation shielded, for example by a combination of low-Z material (for example, 2.5 cm polycetal (POM)) and high-Z material (for example, 1.5 cm lead). There has to be a hole in that shielding with a diameter of a few mm on axis, in order to let on axis-electrons pass into the deflecting permanent magnetic field inside. All other sides of the spectrometer should be shielded additionally, for example with flexible sheet lead with a thickness of the order of a few mm, in order to minimize the background signal. In addition, one can add a holed image plate in front of the spectrometer in order to be able to retrieve information not only on the spectral energy distribution on axis, but also on the divergence of the emitted electrons.

Figure 27 shows a typical setup in a vacuum chamber which can be used to generate and measure electrons from laser-foil interaction using such a spectrometer. The laser beam (red) is incident from the left and is focused by the gold-coated parabolic mirror onto the metallic target foil. Electrons are accelerated and emitted (green) behind the target foil with approximately 30° divergence and are analyzed by the spectrometer. The figure also shows vividly the extremely low spacial footprint of the whole setup. There is no electron beam or hard radiation until the laser beam hits the target foil, and 10 centimeters behind the source, the beam is already being analyzed [110].

As an example, figure 28 shows the electron signal on an image plate in front of the spectrometer (a) and three shots recorded in the detection plane (b). To produce these signals, an 80-fs laser pulse with an intensity of the order of $I \approx 5 \times 10^{19}$ W/cm² and titanium foil targets of a few µm thickness were used.

After post-processing and evaluating the raw signals on the image plates in front of the target and in the detection plane, exponential electron energy spectra can be retrieved, as shown in figures 29 and 30. For the spectra in figure 29, the laser intensity amounted to $I \approx 5 \times 10^{19}$ W/cm², whereas for figure 30, the intensity amounted to $I \approx 2 \times 10^{19}$ W/cm², only. The corresponding drop in temperature as predicted by scalings 69 and 70, for example, is clearly visible.

Furthermore, there is a dependency of target thickness. In figure 29 a), the electron
temperature when using 2 µm thick titanium foils amounted to $T_{\text{eff}} \approx 1.4$ MeV, whereas in b) with 5 µm thick titanium a significantly higher temperature of $T_{\text{eff}} \approx 2.3$ MeV was obtained. Likewise, in case of the lower intensity of $I_L \approx 2 \times 10^{19}$ W/cm$^2$ (see figure 30), different target thickness was used. With 5 µm Ti, the measured temperature was $T_{\text{eff}} \approx 0.6$ MeV, and with 25 µm Ti the effective temperature amounted to $T_{\text{eff}} \approx 0.5$ MeV. It shall be noted that here, the intensity variation was produced by varying the laser energy, not by moving the target out of the focus. Therefore, the focal spot size remained the same at all intensities.

### 7.5 Proton and ion acceleration

The today achievable maximum laser intensities up to $\approx 10^{22}$ W/cm$^2$ are not yet sufficient to move directly protons and ions due to their much higher masses. However, the quasistatic electric fields which are set up by the accelerated electrons during laser-solid interaction can lead to efficient proton and ion acceleration. A natural source of protons and ions can be hydrocarbon contaminants which are typically present, for example, at the back side of target foils. The predominant mechanism which leads to proton and ion acceleration is the so-called "Target Normal Sheath Acceleration" (TNSA), which was discovered by independent groups almost at the same time [132,136–139].

As discussed in section 7.3, and shown in [132,140], for example, an ultrahigh-intensity
7.5 Proton and ion acceleration

Figure 28: Experimentally recorded Image Plate pictures. a) Readout signal on IP in front of the spectrometer shielding with entrance hole. The approximately round, light-grey isocontour created by the electron signal on the Image Plate (IP) is positioned centrally around the entrance hole. b) Electron raw signal on the detection IP behind the shielding. The low energies are located on the left side of the IP, the high on the right hand side.

Figure 29: Measured electron spectra behind target at a) an intensity of about $I_L \approx 5 \times 10^{19} \text{W/cm}^2$ using a 2 $\mu$m thick titanium foil as target, b) at an intensity of $I_L \approx 5 \times 10^{19} \text{W/cm}^2$ and 5 $\mu$m thick titanium.
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Figure 30: Measured electron spectra behind target at a) an intensity of about $I_L \approx 2 \times 10^{19}$ W/cm$^2$ and a) 5 $\mu$m titanium foil, and b) with 25 $\mu$m thick titanium.

laser ($I \lambda^2 > 10^{18}$ Wcm$^{-1}$ $\mu$m$^{-2}$, where $I$ is the intensity and $\lambda$ the laser wavelength) incident on a thin target metal foil with thickness in the 10 micron range, accelerates electrons from the surface of the target in target normal direction. They propagate through the target in a complex way involving cold return currents, and can ionize atoms along their path. Due to the Boltzmann-like electron energy distribution, most of the electrons have low energies of the order of keV. These electrons leaving the target at the rear side build up large quasi-static electric field. The field acts normally to the target surface, has cylindrical symmetry and decreases in transverse direction. Due to the ultrashort duration of the electron bunch and its high charge, close to the axis the quasi-static electric field may reach values of several TV/m. Hence the potential can attain several tens of MeV [139]. Low-Z-ions, mainly protons, present at the backside of the target foil may be accelerated by that field until they compensate the electron charge. The origin of these ions are mostly hydrocarbon contaminates on the target surface [141,142], but they can also be attached deliberately, for example as proton-rich dots [143]. Figure 31 visualizes the basic principle of the TNSA mechanism. Various theoretical models to describe TNSA have been developed, an extensive comparison can be found in reference [144].

Since a preplasma has a substantial influence on the electron acceleration process, as described in section 7.2, and the accelerated electrons are responsible for ion acceleration, it is consequent that there is also an influence of the laser prepulse on proton and ion acceleration [145]. Similarly, the shape of the target (e.g., bent [146] and cone targets [147,148]) has considerable impact on the obtainable proton output, for example as regards focusing [149].

Due to the ultrashort acceleration duration and the fact that the protons are at rest
before acceleration, the transverse emittance of the proton beam has values around $10^{-3}$ mm mrad for 10 MeV protons [150]. In comparison the longitudinal emittance is large, with their spectrum exhibiting a quasi-exponential shape with a distinct cut-off energy [136, 137, 139]. The transverse dimension of the electric field and hence the source size of accelerated protons are much larger than the laser’s focal spot for a plane target [142, 151].

As is generally the case with electrons accelerated from laser-solid interactions, accelerated protons typically do also have exponential energy spectra. In standard TNSA scenarios, the typical maximum proton and ion energies obtained with typical 100 TW-class lasers amount to a few to a few tens of MeV. The maximum laser accelerated proton energies to date amount to approximately 70 MeV, obtained using 80 J laser pulses and specially shaped targets at Los Alamos’ trident laser facility [147]. Scaling laws have been developed which predict increasing proton energies as the laser intensity is increased [152, 153]. Proton energies up to a few hundreds of MeV are predicted using

Figure 31: Target normal sheath acceleration (TNSA) scheme: A TW-laser pulse is focused onto the front side of the target foil, where it generate a blow-off plasma and subsequently accelerates electrons. The electrons penetrate the foil, ionize hydrogen and other atoms at the back surface and set up a Debye sheath, which accelerates the protons.
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upcoming larger lasers with higher intensities. Various different acceleration schemes, such as radiation pressure acceleration (which has its roots in laser propulsion [154]) have been proposed and developed, where the maximum laser energy scales as strong as \( \propto a_0^2 \) with the laser intensity.

While most of the above experiments and concepts yielded proton beams with exponential energy distribution, there are also scenarios where quasimonoenergetic features are observed in laser-overdense interactions are obtained [143, 155, 156], e.g. by using proton-rich dots or mass limited micro-droplets or by using higher wavelength lasers such as CO2-lasers, where gas jet targets might already be overdense to the incident laser light [157, 158]. However, the spectral widths and the energies obtained still leave much room for improvement when compared to state-of-the-art accelerators.

One main motivation for research in laser-driven proton and ion acceleration stems from the idea that laser-plasma produced protons as well as ions might be used effectively for cancer treatment. The sharp Bragg peak of protons and ions due to their stopping powers, allows for depositing large dose in a confined region within the patient. Therefore, such regions can be treated without damaging too much the surrounding tissue. Also, dose escalation for inherently radioresistant tumors was an initial motivation for research in proton and ion therapy research using conventional accelerators [159].

Using laser-plasma-produced protons and ions was first suggested at the beginning of the century [160,161], and a lot of work has been done since then (see [162], for example), and the possibility of using lasers as energetic proton sources for therapy was already mentioned in general reviews on particle therapy [163]. A large number of groups is doing research on this in principle very appealing application. However, among other things, an enormous reproducibility, monochromacity and tunability of the generated proton or ion beams are required for such highly susceptible medical applications such as proton cancer therapy. At the same time, advanced non-laser based accelerators such as specially designed superconducting cyclotrons emerge to the market, which are much more cost-efficient than conventional proton sources and not much larger than a state-of-the-art high-power laser system [164]. Protons and ions are much more easy particles to accelerate in circular accelerators when compared to electrons due to their orders of magnitude lower mass and a therefore substantially lower synchrotron radiation losses. There is substantial doubt that in a field where reliability and controllability are so crucial, laser-produced protons and or ions will ever become competitive [165].

However, there have been also other applications for laser produced protons and ions identified such as proton radiography [166], which has already been demonstrated extensively using lasers [167], for example in the context of tumor imaging [168, 169], laser-driven implosions [170]. Another is using protons as igniter pulse in fast ignition schemes [133].
7.6 Summary

The laser intensity is the most important parameter as regards electron beam output from laser-overdense experiments. At first, it determines to a large extent the temperature and thus, the spectral distribution of the generated electron beam. Well-known scaling laws such as those by Wilks and Beg exist which can be used to determine the temperature of an electron beam. Dependent on which scaling is used, the temperature increases at maximum with the square root of the incident intensity. Typical temperature amount to few MeV, with typical maximum cutoff energies of about 10 MeV.

Second, more recent findings indicate that the divergence of the generated beam is also dependent mainly on the incident laser intensity. Again, the divergence increases as the laser intensity increases. According to available data, expected values amount to about 25° for intensities of a few $\approx 10^{18}$ W/cm$^2$ up to about 55° for intensities of a few $\approx 10^{20}$ W/cm$^2$.

The electrons which are emitted behind the target are accelerated at the front of the target at the vacuum-solid boundary, where the laser is incident. Part of the laser energy is converted into hot electrons, and the majority of the laser energy is reflected near the critical density. The accelerated electron current which propagates into the solid exceeds the Alfvén current, and within the target therefore cold return currents are generated. Therefore the target material (typically metals such as aluminum), thickness (typically of the order of 10 microns) and the preplasma (typical scale length of $[d(\log N_e)/dx]^{-1} \approx 200$ nm) do also have substantial influence on the generated electron beam parameters.

Hydrodynamic simulation tools can be used to determine the preplasma situation after arrival of a prepulse just before incidence of the main laser pulse. The main laser-target interaction can then be modeled using particle-in-cell simulation tools using input on the preplasma situation obtained with the hydrodynamic approach.

Data on the charge of emitted electrons is scarce. Typical expected and measured values are of the order of tens of nC. The laser intensity on target can be varied either by changing the laser pulse energy or by moving the target out of focus. While in the first case, the laser spot size remains constant, in the latter case the spot size increases substantially, which does probably influence the total charge accelerated.

Measurement of accelerated electrons is possible by using electron-sensitive phosphor screens or image plates, for example to measure the divergence. When combined with compact electron spectrometers based on permanent magnets, the energy spectrum can be retrieved. Absolute calibration of image plates allows to determine the charge.

A wide range of laser systems (from sub-TW to PW powers) has already been used to produce energetic electrons from interaction with solids. A lot of that research is in the context of the fast ignition scheme for inertial confinement fusion, and by proton and ion acceleration schemes (see section 7.5). Proton and ions produced by laser-plasma-interaction can therefore also be useful for a wide range of applications. As is the case for electrons, the laser intensity is again of paramount importance as regards the proton and
ion output. Higher energy electrons with much smaller divergence and energy spread can be produced from laser-underdense interaction (compare section 8).


8 Laser-underdense interaction

8.1 Basic principles

In contrast to the overdense case discussed in section 7, electromagnetic waves can propagate through underdense media with a phase velocity of $v_{ph} = c/n_p$ and group velocity

$$v_{g,L} = c \eta_p = c \sqrt{1 - \left(\frac{\omega_p}{\omega_L}\right)^2}$$

(86)

where $\eta_p$ is the index of refraction of the plasma. A focused laser high-power laser pulse with intensities of the order of $\approx 10^{18}$ W/cm$^2$ or higher not only quasi-instantaneously ionizes gaseous media, but also expels off axis, which are then re-attracted by the quasi-static ion background. The dynamics of this re-attraction is determined by the local plasma frequency $\omega_p$. Parameters of paramount importance are the plasma wavelength $\lambda_p$ and the laser pulse duration $\tau$. If the laser pulse duration is smaller than the half of the plasma wavelength, $c\tau < \lambda_p/2$, the conditions are right to form a strong longitudinal plasma wave and an electron-free plasma blowout [171] which trails the driving laser pulse with a phase velocity of

$$v_{ph,pw} = v_{g,L} = c \sqrt{1 - \left(\frac{\omega_p}{\omega_L}\right)^2}$$

(87)

The difference to the laser-overdense case is that the accelerating plasma structure is not stationary, but moves through the plasma with the speed of light (in plasma), and can therefore be regarded as a relativistically moving plasma cavity. This electric wakefield can then be used to accelerate electrons to extremely high energies.

In the linear regime ($a_0^2 \ll 1$) the plasma wave generation can be described based on the linearized equation of motion, the continuity and Poisson equation [172–174]. This leads to

$$\left(\frac{\delta^2}{\delta t^2} + \omega_p^2\right) \frac{\delta n}{n_0} = -\omega_p^2 c^2 \nabla^2 \frac{a^2}{2}$$

(88)

The ponderomotive term on the right hand side of the equation is a result of the laser amplitude $a$, whereas the restoring, electrostatic force is the result of the plasma density perturbation $\delta n/n_0$. Equation can now be solved by (88)

$$\frac{\delta n}{n_0} = \frac{c^2}{\omega_p} \int_0^t dt' \sin \omega_p(t-t') \nabla^2 \frac{a^2}{2}$$

(89)

which leads, following [174], to an axial wakefield

$$E(r,t) = -\frac{m_e c^2 \omega_p}{e} \int_0^t dt' \sin \omega_p(t-t') \nabla^2 \frac{a^2(r,t)}{2}$$

(90)
For example, for a linearly polarized laser pulse of length $c\tau \approx \lambda_p$ focused to a spot of radius $r_s$ and an intensity profile $a^2 = a_0^2 \exp(-2r^2/r_s^2) \sin^2(\pi(z - ct)/\lambda_p)$ one gets a density perturbation of

$$\frac{\delta n}{n_0} = -\frac{\pi}{16} a_0^2 \left[ 1 + \frac{8}{k_p r_s^2} \left( 1 - \frac{2r^2}{r_s^2} \right) \right] \exp \left( -\frac{2r^2}{r_s^2} \right) \sin(k_p(z - ct))$$  \hspace{1cm} (91)

and an electric field of

$$\frac{E}{E_{wb,\text{klass}}} = -\frac{\pi}{16} a_0^2 \exp \left( -\frac{2r^2}{r_s^2} \right) \cos(k_p(z - ct))$$  \hspace{1cm} (92)

where $E_{wb,\text{klass}}$ is the classical wave breaking limit defined by

$$E_{wb,\text{rel}} = E_{wb,\text{klass}} \sqrt{2(\gamma_{ph,pw} - 1)}$$  \hspace{1cm} (93)

and $\gamma_{ph,pw} = (1 - \beta_{ph,pw}^2)^{-1/2}$ is the Lorentz factor of the longitudinal plasma wave. In the linear, one-dimensional case the maximum electric field which can be transported and sustained by a plasma wave without breaking amounts to

$$E_{wb,\text{klass}} = c\omega_p m_e/e$$  \hspace{1cm} (94)

or as an engineering formula [174]

$$E_{wb,\text{klass}} [\text{V/m}] \simeq 96 \sqrt{n_0 [\text{cm}^{-3}]}$$  \hspace{1cm} (95)

As an example, the classical wavebreaking limit amounts to $E_{wb,\text{klass}} \simeq 100 \text{GV/m}$ for a gas density of $10^{18} \text{cm}^{-3}$. This is three or four orders of magnitude higher than in conventional accelerator systems. At the same time, there are transversal electric fields of the same order of magnitude in a plasma wave which focus electrons – another highly desirable characteristic of plasma waves.

Wave-breaking can lead to injection of background plasma electrons into the accelerating and focusing phase of the plasma wave, where they can be accelerated to high energies on very short distance. It shall be noted that wave-breaking is a highly complex phenomena in case of a warm, three-dimensional plasma, for example because in a warm, thermal plasma there is always a small (exponentially decreasing) fraction of plasma electrons with velocities above the phase velocity $v_{ph}$ of the plasma wave. Wave-breaking is discussed in detail for the 1D-case in [175].

### 8.2 Energy gain in plasma waves

The maximum energy $W^{\text{max}}$ which can be obtained by electrons in a plasma wave is given by the integral of the longitudinal $E$-field $E(z)$ over the propagation distance $L$

$$W^{\text{max}} = e \int_0^L E(z) dz$$  \hspace{1cm} (96)
8.2 Energy gain in plasma waves

Figure 32: a) Schematic visualization of dephasing in a longitudinal (linear) plasma wave. While single electrons can propagate in the forward direction with a velocity $v_e \simeq c$ close to the vacuum speed of light, the phase velocity of the plasma wave $v_{ph,pw} < c$ is significantly lower than the speed of light. Therefore, the fast electrons move forward with respect to the plasma blowout cavity and sooner or later leave the accelerating phase of the plasma wave and cannot gain energy anymore. b) Density perturbation (right y-axis) and corresponding electric field strength (left y-axis) of such a linear plasma wave.

The electric field $E_z$ experienced by an injected electron is not constant, and the acceleration distance $d_z$ is limited. The three fundamentally limiting factors which limit the acceleration distance and the maximum energy gain $W_{max}$ are

- laser depletion
- dephasing
- diffraction of the laser pulse

8.2.1 Laser energy pump depletion

Energy is transferred from the laser pulse to the plasma in order to excite the plasma wave. The distance $L_{pd}$ after which the total laser pulse energy would be lost to the plasma can be derived by equating the plasma wave energy with the laser pulse energy [176]. The energy of the plasma wave $W_{pw}$ is the product of volume $V_{pw} = \pi r_0^2 L_{pd}$ and energy density $U_{pw} = \epsilon_0 E_z^2/2$, and for the laser pulse with an electric field $E_L$ we have a laser pulse volume $V_L = \pi r_0^2 c \tau$ and $U_L = \epsilon_0 E_L^2$. The electric field of the laser pulse amounts to $E_L = a_0 m_e \omega_0 / e$, scaling linearly with $a_0$, while the longitudinal wakefield scales as $a_0^2$ (compare equation (92) and reference [177] and [178]). In case of a rectangle
pulse with $c \tau = \lambda_p / 2$ it can be given as $E_z = a_0^2 m_e c \omega_p / e$. Equating leads to

$$W_L = V_L U_L = V_{pw} U_{pw} = W_{pw}$$

$$\pi r_0^2 c \varepsilon_0 E^2_L = \pi r_0^2 L_{pd} \varepsilon_0 E^2_z / 2$$

$$\lambda_p E^2_L = L_{pd} E^2_z$$

$$\lambda_p \omega^2_L = L_{pd} \omega^2_p a_0^2$$

and

$$L_{pd} = \frac{\omega^2_L \lambda_p}{\omega^2_p a_0^2}$$

for the non-relativisite case $a_0 \ll 1$, respectively. This result means that the depletion length does not actually decrease as laser intensity increases in the linear case, because the longitudinal wakefield increases faster than the electric laser field.

When there are linear plasma waves of type $E_z = E_{max} \sin[\omega_p (z/v_{ph} - t)]$, one has $E_{max} \ll E_{wb,\text{lass}}$, and the corresponding plasma wave has a comparatively small amplitude $\delta n / n_0$ and a structure which can be described by a sine function such as given in figure 32. At larger fields $E_{max} \rightarrow E_{wb,\text{lass}}$, however, one reaches the nonlinear regime, where the electric field steepens and gets a saw-tooth shape and the plasma density has sharp maxima [179]. The electron density variations can amount up to 100% and lead to nearly complete electron explosion and formation of electron-free cavities. An increase of the plasma wave amplitude does also lead to higher electron oscillation velocities and a significant increase of plasma wavelength. this nonlinear plasma wavelength $\lambda_p^{NL}$ amounts to [178]

$$\lambda_p^{NL} = \frac{2 \pi}{E_{max} / E_{wb,\text{lass}}}$$

Since for example for a circular polarized laser pulse with rectangle profile the maximum electric field is $E_{max} = a_0^2 (1 + a_0^2)^{-1/2}$, with the nonlinear plasma wavelength $\lambda_p^{NL}$ in contrast to the linear case (98) we have a nonlinear depletion length $L_{pd}^{NL}$ of [178]

$$L_{pd}^{NL} = \frac{\omega^2_L \lambda_p a_0}{\omega^2_p 3\pi}$$

Therefore, in this case the depletion length increases as intensity increases.

### 8.2.2 Dephasing

Even stronger limiting factors than laser pulse depletion are dephasing and diffraction. The massive kinetic energy gain of electrons in the GV/m scale electric fields of the plasma wave means, that these electrons quickly approach the velocity of light in vacuum $c$. However, since the phase velocity of the plasma wave is smaller than the speed of
light in plasma \((v_{ph,pw} = v_{g,L} < c\), see equation (87)), the electrons in the accelerating phase of the plasma wave tend to outrun the laser driver and reach the second phase of the plasma wave, where the electric field is decelerating the electrons. Figure 33 shows the propagation speed difference between relativistic electrons in dependence to their kinetic energy (left/lower axis), and the laser pulse in dependence of the plasma density (right/upper axis).

![Figure 33: Comparison of propagation velocities of electrons and a laser pulse in plasma.](image)

The left \(y\)-axis shows the electron \(\beta = v/c\) in dependence on its energy, and the right \(y\)-axis denotes the laser pulse group velocity \(\beta_{g,L} = \left(1 - n_e/n_c\right)^{1/2}\).

The velocity difference \(v_e - v_{ph,pw} > 1\) between electron and plasma wave can now be used to calculate the dephasing or detuning time which is needed for the half plasma wave \(\lambda_p/2\), which yields

\[
t_d = \frac{\lambda_p}{2(v_e - v_{ph,pw})} = \frac{\lambda_p}{2c(\beta_e - \beta_{ph,pw})}.
\]  

The corresponding dephasing length \(L_d\) is an experimentally better accessible parameter. To derive that length, one sets \(v_e \approx c\) and \(\beta_e \approx 1\), respectively. Since the plasma is undercritical and therefore \((\omega_p/\omega_L)^2 < 1\), it is allowed to develop the wake’s phase
velocity \( v_{\text{ph, pw}} \) from equation (87) and gets \( \beta_{\text{ph, pw}} \simeq 1 - 1/2(\omega_p/\omega_L)^2 \), yielding

\[
L_d = \frac{\lambda_p}{2(\beta_e - \beta_{\text{ph, pw}})} \simeq \frac{\lambda_p}{2} \gamma_{\text{ph, pw}} \simeq \lambda_p \left( \frac{\omega_L}{\omega_p} \right)^2 = \lambda_p \frac{n_e}{n_e} \tag{102}
\]

Here, \( \gamma_{\text{ph, pw}} = (1 - (v_{\text{ph, pw}}/c)^2)^{-1/2} \) is the Lorentz factor associated with the phase velocity of the plasma wave.

Based on this dephasing length \( L_d \), the maximum energy gain \( W_{\text{max}}^{\text{ph}} \) of an electron according to equation (96) and assuming a constant field \( E_{\text{wb, klass}} \) (see equation 94) is then

\[
W_{\text{max}} = e E_{\text{wb, klass}} L_d = 2\pi \left( \frac{\omega_L}{\omega_p} \right)^2 m_e c^2 \tag{103}
\]

Decreasing the plasma density and therefore \( \omega_p \) therefore leads to a larger potential energy gain: the group velocity \( v_{g,L} \) of the laser pulse and therefore the phase velocity of the laser-driven wakefield is closer to the speed of light in vacuum \( c \), which means that electrons can interact longer with the accelerating wakefield.

However, it shall be noted that decreasing the plasma density to achieve larger energy gains on the other hand further increases the problems associated with diffraction of the laser pulse, see next section.

In order to extend the dephasing to the nonlinear, highly relativistic case \( (\alpha_0^2 \gg 1) \), the nonlinear plasma wavelength \( \lambda_p^{NL} \) (see equation (99)) must be used. For a matched-duration laser pulse with rectangular profile the wakefield is maximum \( E_{\text{max}}/E_{\text{wb, klass}} = \alpha_0^2 \) \cite{178}. Again, this is valid for a circular polarized laser pulse, while one gets \( E_{\text{max}}/E_{\text{wb, klass}} = \alpha_0^2/2 \) for linear polarization \cite{178}. Therefore the nonlinear dephasing length for a circularly polarized pulse is

\[
L_{\text{d}}^{NL} \approx \lambda_p \left( \frac{\omega_L}{\omega_p} \right)^2 \frac{2}{\pi} \frac{E_{\text{max}}}{E_{\text{wb, klass}}} \simeq \lambda_p \left( \frac{\omega_L}{\omega_p} \right)^2 \frac{2}{\pi} \frac{a_0^2}{\pi} \tag{104}
\]

Table 1 summarizes the expressions for dephasing and depletion lengths for the non-relativistic and the relativistic case.

<table>
<thead>
<tr>
<th></th>
<th>non-relativistic ( (\alpha_0^2 \ll 1) )</th>
<th>relativistic ( (\alpha_0^2 \gg 1) )</th>
</tr>
</thead>
<tbody>
<tr>
<td>depletion</td>
<td>( L_{\text{pd}} = \frac{n_e}{n_e} \lambda_p \frac{1}{\alpha_0^2} )</td>
<td>( L_{\text{pd}}^{NL} = \frac{n_e}{n_e} \lambda_p \frac{a_0}{\pi} )</td>
</tr>
<tr>
<td>dephasing</td>
<td>( L_d = \frac{n_e}{n_e} \lambda_p )</td>
<td>( L_d^{NL} = \frac{n_e}{n_e} \lambda_p 2\frac{a_0^2}{\pi} )</td>
</tr>
</tbody>
</table>

Table 1: Formulas for depletion and dephasing lengths in non-relativistic and relativistic scenarios.

In addition to the longitudinal electric field, in the linear, three-dimensional case in addition there is also a radial component, which exerts a strong focusing force. This
transversal $E$-field (and a transversal $B$-field) is the result of the Panofsky-Wenzel theorem [180, 181], which describes the conjunction of the longitudinal $E$-field with the transversal fields via $\partial E_z/\partial r = \partial (E_r - B_\theta)/\partial (z - ct)$ [178].

Several strategies have been developed to overcome dephasing in laser-plasma-accelerators, for example the surfatron [182–184]. Here, the maximum energy gain as given by equation (103) shall be increased by avoiding the dephasing by adding a constant magnetic field which is perpendicular to the transverse electric field. This magnetic field of the order of a few to a few tens of Tesla would produce an additional $v \times B$-Drift which could compensate for the phase delay, similar to a surfer who surfs in an angle to the wave in order to get a longer ride. This idea is hard to realize, but nevertheless it was used to contribute to explain the synchrotron radiation observed during supernova phenomena, which assumedly is generated by highly relativistic electrons [185, 186].

Another strategy is density tapering [187], where a plasma density upramp would lead to a contracting blowout, which could also be used to keep an accelerated bunch at a constant phase or longitudinal electric field within the blowout for an extended distance.

### 8.2.3 Diffraction

Another highly important limiting factor is diffraction of the laser pulse. The linear propagation of a conventional Gaussian laser pulse can be described via the Rayleigh length $Z_R = \pi \omega_0^2/\lambda_L$, where $\lambda_L$ and $\omega_0$ are the laser wavelength and laser spot waist size, respectively (see above). The opening angle of a Gaussian beam far behind the focus is $\alpha = dR/dz = \omega_0/Z_R = \lambda_L/(\pi \omega_0)$. Practically, this leads to a quick decrease of light intensity and would soon lead to an inability to drive plasma waves after short acceleration distance. Figure 34 a) shows schematically the diffraction of a Gaussian beam in vacuum and in b) reproduced based on a 3D-PIC-Simulation using PLASMA SIMULATION CODE. Here, an 80-fs long Ti:sapph laser pulse was injected into the simulation box from the left, focused to a pulse width (FWHM) of 2.5 $\mu$m. The 2D-Plot of the electric field in the polarization plane shows clearly the broadening of the laser pulse due to wave front curvature and diffraction.

As an example, the Rayleigh length of a Ti:sapph laser pulse focused to a spot width of 5 microns would result in a Rayleigh length of $\approx$ 150 microns, only.

**Self-focusing**

Fortunately, intense laser-plasma interaction can lead to self-focusing of a laser pulse, and therefore to an effect counteracting the natural diffraction of a laser pulse. Generally speaking, this effect sets in if there is a transversal plasma density profile which has a lower electron density on the laser axis when compared to away from the laser axis. Such a radial plasma density perturbation $\partial n_e/\partial r > 0$ is leading to a radial decrease
Figure 34: a) Schematic drawing on the broadening of a Gaussian beam due to diffraction. b) PIC-simulation with psc showing diffraction in vacuum. An 80-fs laser pulse focused to a FHWM width radius of 2.5 µm and an intensity of $I = 2 \cdot 10^{18}$ W/cm$^2$ has been sent into the simulation box from the left. The laser E-field in the polarization plane is plotted, which shows the bending of the local wave fronts due to diffraction.

of the index of refraction $\eta = (1 - \omega^2 / \omega^2_p)^{1/2}$ and consequently to a radial increase of the phase velocity $c/\eta$. The resulting bending of wavefronts leads to self-focusing of the laser pulse. This is shown schematically in figure 35 a), and is complemented in b) based on a 2D-PIC-simulation 2D-PIC-Simulation with psc [188] for a laser pulse intensity of $I = 10^{19}$ W/cm$^2$, a width (FWHM) of 5 µm and a gas density of $n_e = 10^{20}$ cm$^{-3}$.

Such a self-focusing can have different contributing factors, dependent of duration and intensity of the laser pulse and the relevance of collisions [189, 190]. In all variations, however, the self-focusing is ultimately based on the (often parabolic) radial decrease of the index of refraction $\partial \eta / \partial r < 0$ and in this respect resemble convex lenses known from optics.

Before the development of chirped pulse amplification (CPA) and the availability of ultrashort laser pulses with relativistic intensities self-focusing was often produced by variations of the plasma density due to ion movement. However, with ultrashort, relativistic laser pulses there are additional effects which are generated by the movement of electrons in the laser field alone, when the ions are quasistatic. One can break down the individual contributions which may be responsible for changing the index of refraction into

**Thermal self-focusing due to collisions (Joule heating).** In plasmas which are not collision-free collisional absorption can lead to local heating of the plasma. Due to the radially decreasing laser intensity $\partial I / \partial r < 0$ the thermal pressure $n_0 k_B T_e$ is larger on axis than further outside, leading to an increased hydrodynamic expansion and in consequence to a radially increasing plasma density $\partial n_e / \partial r > 0$ [191].

**Ponderomotive self-focusing.** At higher intensities, electrons are being displaced away from the axis ponderomotively. Again, this leads to a radial change of plasma
electron density and therefore to self-focusing. Dependent on laser pulse duration and intensity one may differentiate between the following two cases:

– the ions have enough time during the laser pulse passage to follow the electrons. Thermal self-focusing as well as this type of self-focusing (or combinations of both [192]) happen on the timescale of the ion movement, i.e. with a time constant of the reciprocal ion plasma frequency (and not with the much higher electron plasma frequency) \( \tau = 1/\omega_{\text{pl}} = (m_i e_0/(n_i e^2))^{1/2} \).

– only electrons are displaced from axis during the passage of especially short an intense laser pulses, leaving quasistatic ions behind.

Relativistic self-focusing [174, 193–197]. As is the case in nonlinear optics, the index of refraction has an intensity-dependent term. When dealing with ultrashort, relativistic laser pulses \( \tau \ll 1 \text{ ns} \) interacting with plasma \textit{in statu nascendi} there may be relativistic mass increase and therefore a decrease of the plasma frequency due to oscillation movement of electrons in the electromagnetic field of the laser pulse and generation of plasma waves with especially large amplitude \( n/n_0 \). Therefore (and under the assumption that \( \omega_p^2/\omega_L^2 \ll 1 \), which allows to use the Taylor series as above) the resulting nonlinear index of refraction is

\[
\eta_{\text{rel}} \simeq 1 - \frac{\omega_p^2}{2 \gamma(r) \omega_L^2} \tag{105}
\]

and therefore is dependent on the plasma density \( n_e(r) \) and also of the relativistic Lorentz factor \( \gamma(r) \). The dominating contribution to the motion of an electron comes from the oscillation in the laser field \( p_{\perp} = m e a \). Therefore one can express the relativistic Lorentz factor as \( \gamma \simeq \gamma_{\perp} \simeq (1 + a^2/2)^{1/2} \). Using this for the relativistic index of refraction one gets

\[
\eta_{\text{rel}} \simeq 1 - \frac{\omega_p^2}{2 \omega_L^2} (1 + a^2(r)/2)^{-1/2} \tag{106}
\]

Assuming that relativistic self-focusing sets in early \( (a_0^2 < 1) \), the index of refraction can be approximated as

\[
\eta_{\text{rel}} \simeq 1 - \frac{\omega_p^2}{2 \omega_L^2} \left( 1 - \frac{a^2(r)}{4} \right) \tag{107}
\]

Using the laser pulse envelope \( a(r) = a_0 \exp(-r^2/(2r_0^2)) \) one can express the propagation of phase fronts with radially decreasing phase velocities \( v_{\text{ph}}(r) = c/\eta \) as

\[
\frac{v_{\text{ph}}(r)}{c} = \frac{1}{\eta} \simeq 1 + \frac{\omega_p^2}{2 \omega_L^2} \left( 1 - a_0^2 \exp\left(-r^2/(2r_0^2)\right) / 4 \right) \tag{108}
\]
Figure 35: a) Schematic visualization of self-focusing. b) 2D-PIC-simulation result obtained with PSC on relativistic self-focusing for a laser pulse with duration 80 fs, for an intensity of $I = 10^{19}$ W/cm$^2$, a vacuum waist width of 5 µm (FWHM) and a gas density of $n_e = 10^{20}$ cm$^{-3}$. The square of the $E$-field of the laser pulse in the polarization plane is shown. The upper snapshot was taken about 60 fs after the simulation entered the simulation box from the left, the bottom snapshot shows the situation 20 fs later. The relativistic self-focusing increases substantially as the propagation distances increases.

and obtains the phase velocity difference

$$\frac{\Delta v_{ph}(r)}{c} = \frac{a_0^2 \omega_p^2}{8 \omega_L^2} \exp \left(-r^2/(2r_0^2)\right)$$  \hspace{1cm} (109)

This curved phase front manifests in a maximum delay between a phase front far away from axis ($r = \infty$ and $a = 0$) and a phase front on axis ($r = 0$) as $\Delta L = |c \Delta v_{ph}(r)|_{max} t = |\Delta v_{ph}/c|_{max} Z = \alpha R$. Geometric considerations (see figure 34 a)) for small angles yield the relations $\alpha \approx \sin \alpha \approx \tan \alpha = R/Z = \Delta L/R$, and one obtains $\alpha^2 = (\omega_p a_0)^2/(8 \omega_L^2)$ quantitatively as the angle which would be the effect of relativistic self-focusing in absence of any diffraction [196, 197].

**Quantification of the limit for relativistic self-focusing**

Since defocusing due to diffraction is opposed to self-focusing, one can come up with a balance equation. In the following, this is done using relativistic self-focusing. Expressing the defocusing due to diffraction as seen above and in figure 34 using the angle $\alpha = dR/dZ = r_0/Z_R = c/(\omega_p a_0)$ and equating this with the counteracting focusing angle

$$\frac{c^2}{\omega_L^2 \omega_0^2} = \frac{(\omega_p a_0)^2}{8 \omega_L^2}$$  \hspace{1cm} (110)

leads to

$$\frac{a_0^2 \omega_0^2 \omega_p^2}{c^2}$$  \hspace{1cm} (111)
Using this equation one can determine when relativistic self-focusing is stronger than diffraction. The product of intensity and area $a_0^2\omega_0^2$ is proportional to the laser pulse power, which results in a power limit. The exact limit is

$$P_L > P_{c,rsf} = \frac{2cm^2e^4\omega_0^2}{\omega_p^2} \approx 17\frac{\omega_0^2}{\omega_p^2} [\text{GW}] = 17\frac{n_e}{n_e} [\text{GW}] \quad (112)$$

The laser pulse power $P_L$ has to be higher than the critical density $P_{c,rsf}$ in order to yield net focusing.

Therefore the relativistic self-focusing limit is dependent exclusively of plasma density and laser wavelength, and not of focusing of the laser beam to a special spot size or its vacuum intensity.

Consequently, the limitation of acceleration length due to diffraction can be overcome due to relativistic self-focusing. The laser pulse and the plasma wave driven by the laser can then propagate over many Rayleigh lengths, and electrons can be accelerated to large energies.

Figure 35 b) is the result of a 2D-PIC-simulation made using plasma simulation code for a plasma density of $n_e = 10^{20}\text{ cm}^{-3}$ and a laser pulse with a duration of 80 fs at an intensity of $I = 10^{19}\text{ W/cm}^2$ and an initial width (FWHM) of 5 $\mu$m. The squared $E$-field of the laser pulse in the polarization plane visualizes transversal self-focusing of the laser pulse, which enters the simulation box from the left hand side. It can be seen clearly that relativistic self-focusing dominates clearly over diffraction. Shown are two snapshots with a distance of about 20 fs in order to give an impression on the development of the self-focusing.

**Ionization defocusing**

In case of a laser of sufficient intensity and the propagation through a non- or only partly ionized gas there is in addition an ionization defocusing effect. Here, the release of electrons due to additional ionization of higher ionization levels leads to a growth of plasma density on axis $\partial n_e/\partial r < 0$, which counteracts self-focusing. It may happen that a locally decreasing index of diffraction leads to defocusing, but then self-focusing overcompensates diffraction and gets dominant, the laser pulse is focused and re-enters the regime where ionization defocusing sets in and so forth. In such a scenario, periodic ionization sparks can be formed ([198] and references therein).

Next to self-focusing effects, it is possible to generate a suitable plasma profile which allows for extended propagation actively. For example, a laser can be used to preform a plasma channel with a transverse parabolic density profile [49,199], or triggered discharges in capillaries can be used to produce plasma waveguides [200–202]. Figure 36 shows schematically how a laser pulse is guided over cm-scale distance in a capillary plasma waveguide. The radial plasma density profile is indicated by the b/w color coding.
capillary waveguide with parabolic plasma density

Figure 36: Schematic visualization of guiding of a laser pulse in a capillary. A properly shaped radial plasma density profile (b/w color coding) within the capillary can extend the guiding over cm-scale distance. The longitudinal plasma waves can therefore also be driven over a distance substantially longer then in the vacuum diffraction case and therefore may enable electron acceleration to substantially higher energies.

As a final note, it shall be emphasized that increasing the dephasing distance and the energy gain by working at lower plasma densities and/or by density tapering in turn increases the problem of diffraction, because then the laser pulse must be guided over longer distances.

8.3 Injection, trapping and acceleration

However, controlling the injection of electrons into the proper phase of the blowout and trapping is a complex task. Laser-plasma-accelerators as we know them today were first brought into discussion about 30 years ago in the seminal publication “Laser Electron Accelerator” [203]. The was in the pre-CPA era, and when more and more high-power laser systems started to get available at the end of the 1980s, this gave a boost to the whole field and more and more experiments could be carried through. For example, it was shown that externally generated and injected electron beams can be further accelerated with ultrahigh-gradient in a laser-excited relativistic electron plasma wave [204]. In contrast, one further especially appealing feature of laser-plasma acceleration is the ability to use electrons from the background plasma wave itself to constitute an accelerated beam, negating the need for an externally generated/injected beam. Early experiments [205–208] on this were characterized by exponential electron spectra with cutoff energies up to 200 MeV [208]. In 2002, another milestone was reached when it was shown based on particle-in-cell simulations, that self-injected electrons can be accelerated to quasi-monoenergetic electrons [209]. Figure 37 a) shows schematically, how electrons are injected from the vertex of the blowout and form a stem of accelerating electrons inside the bubble. The (highly nonlinear) longitudinal electric field distribution inside the bubble is also indicated. The turning point of the accelerating field is approximately in the middle of the blowout. Electrons which previously have been ejected off axis by the driving laser pulse may be injected at the vertex of the bubble, on trajectories as indicated by the green arrow [210]. The injection process is highly sensitive on small
variations of the bubble shape, for example, which in turn is dependent not only on the background plasma or gas density, but also on the driving laser pulse power, intensity and size. Modeling and predicting self-injection and trapping is still an issue of intense research and discussion today [211–213]. In a nutshell, an electron is trapped when its forward energy during the injection process quickly enough increases to the phase velocity of the plasma wake, so that the electron can co-propagate. Fortunately, due to the low mass of the electron and the high plasma electron density (which reduces the laser pulse group velocity), this is already the case at electron energies of a few MeV.


Figure 37: Bubble acceleration scheme. Schematic summary of the processes which lead to the giant electromagnetic fields which are exploited by laser plasma acceleration.

Once injected and trapped, the electrons are accelerated to larger and larger energies and corresponding velocities which are soon larger than the phase velocity of the driving plasma wake. Therefore the electrons move forward within the bubble, which leads to the highly undesirable dephasing issue. On the other hand, this dephasing also has the highly advantageous effect of forming monoenergetic features. Consider continuous injection of electrons. Electrons which have already propagated beyond the dephasing point, where the longitudinal electric field is zero, have already gained the dephasing energy. However, they are then decelerated as they enter the second half of the bubble, where the electric field is decelerating. At the same time, electrons which have been injected later are quickly accelerated to large energies, since the electric field is largest.
Laser-underdense interaction

at the vertex of the bubble, and therefore may have the same energy as those energies which have already traveled past the dephasing position. Therefore an energy close to the dephasing energy acts as an attractor, which contributes to forming monoenergetic spikes.

For the first time, quasimonoeenergetic electrons with energies of the order of tens of MeV were reported in 2004 [47–50]. Nature magazine took these results to design the corresponding issue cover, which signals that these experiments were seen as a paradigm change. Since then, many groups have demonstrated the production of monoenergetic electrons with a variety of laser systems [40, 214–219] (to name just a few). Capillary-guided LWFA experiments have also been carried out in order to increase the electron energy to the GeV level [201, 202].

Various injection techniques have been proposed and demonstrated, such as colliding laser pulses [220–223], plasma density transition [224–229] and multi-ionization level injection [230–235]. These injection schemes, or combinations of these schemes, are currently being examined by various groups to enhance the control of injection and acceleration, and to improve beam parameters.

One of the most intriguing characteristics of monoenergetic electron bunches produced by LWFA are their extremely small beam sizes, which can be of the order of microns or even lower. This is a direct consequence of the small accelerating structures, i.e. bubble sizes of the order of few tens of microns. The bunch duration can be as low as a few fs [236–239], and the transverse size can be similarly small. This paves the way for several important applications such as light radiation sources [240]. For example, these electron bunches can be used for generation of bright synchrotron [241] and soft x-ray [242] generation during their passage through undulators based on permanent magnets. Electromagnetic wiggler fields [243] can alternatively be also produced by plasmas [244]. The betatron motion in such plasma wigglers can generate x-ray radiation [245–247] and even $\gamma$-radiation [248]. These and various other applications are reviewed in [249], for example.

Plasma waves can be driven either by intense laser pulses or by compact electron pulses. The compactness of electrons bunches is highly desirable for the latter scenarios, since then the electron pulse may have the ability to self-ionize gases on virtue of its large electric field, negating the need to preionize the medium. The radial electric field $E_r$ of a Gaussian electron bunch is highly dependent on bunch duration $\sigma_z$ and transverse size $\sigma_r$, scaling as

$$E_r(r) \approx \frac{Q}{(2\pi)^{3/2} \sigma_z \epsilon_0 r} \left( 1 - \exp\left( -\frac{r^2}{2 \sigma_r^2} \right) \right)$$

Hybrid systems, which utilize the compactness of electron bunches to ionize media and to drive intense electron waves, have therefore been proposed and examined recently [250–252]. They have the potential to enable unprecedented electron bunch emittance
\( \epsilon_n \), and therefore unprecedented brightness \( B \propto \epsilon_n^2 \). Figure 38 shows the results of a Particle-In-Cell simulation using the code OOPIC [253]. Here, a driver/witness scenario is modeled, where one electron bunch acts as a driver and a second one gains energy as a witness bunch in the plasma wave [250].

![Figure 38: Driver/witness type high-density plasma electron acceleration. The driver electron bunch generates the plasma wave, and the witness bunch is accelerated in the plasma wake. a) shows the electron density distribution, b) the transverse electric field \( E_r \), and c) the longitudinal field \( E_z \) distribution based on a PIC-simulation with OOPIC [253], are exploited by laser plasma acceleration.](image)

Upcoming PW-class laser systems are expected to enable laser-plasma electron acceleration up to energies beyond the 10 GeV level [254].
9 Proof-of-Concept: Space Radiation Testing with Laser-Plasma-Accelerators

It was proposed in Hidding patent, Hidding NIMA [13] to establish laser-plasma-accelerators as novel radiation sources to test electronics, and in particular for space radiation hard-ness testing. The most compelling reasons for this are that laser-plasma-accelerators are becoming more and more compact (table-top) and cost-effective on the one hand, and one the other are extremely versatile, being able to produce electrons, protons and ions as well as (hard) photons with little no to change in the experimental setup, even to produce electrons, protons and ions at the same time. The rapid acceleration in the giant plasma fields tends to produce broadband spectral flux – which is a fundamental characteristics shared with most types of space radiation.

However, the different acceleration environment demands for the development of novel test procedures. Design considerations for laser-plasma-accelerators for space radiation testing have been published in [5] as part of this project. The production of killer electrons which are present in the radiation belts of Earth was chosen as an ideal first step and proof-of-concept case.

The radiation belts of Earth, the van Allen belts, consist of an inner and an outer radiation belt, separated by a ”safe zone”. While the inner radiation belt is dominated by energetic protons, the outer radiation belt is dominated by electrons with energies up to ∼ 10 MeV, sometimes dubbed ”killer electrons” due to their hazardousness for electronics as well as biological systems. The outer radiation belt typically extends from an altitude of about 3 to 10 Earth radii \( R_E \), corresponding to some 13,000 to 60,000 km, with the highest electron flux usually encountered between 4-5 \( R_E \). Understanding of the nature of the acceleration mechanisms, which do involve plasma waves, is subject to ongoing research [255–257]. Same holds for the mapping and prediction of particle flux in the radiation belts, which is also a highly active field of research since the discovery of the van Allen belts [258]. The electron flux within the belts, as well as the extension of the belts, can vary substantially and is connected to solar activity. Figure 39 a) gives an overview on the electron flux predicted by NASA’s standard model AE8\(_\text{max} \) [4] during solar maximum activity for Earth radii ranging from \( L = 1 – 10 \), a value \( B/B_0=1 \) in the energy range of \( E = 0.1 – 7 \) MeV. As can be seen, the electron flux is especially strong at Earth distances which are used by satellite systems near MEO (medium Earth orbit) as important as GPS and upcoming Galileo, and where the geosynchronous orbit (GEO) is located. The same data are visualized in a 3D color plot in figure 39 b). The z-axis is scaled logarithmically, and the grid lines indicate clearly, that mostly the flux energy distribution follows an exponential decay for different \( L \)-values.

This means, that the energy distribution of the van Allen belt electron flux for certain \( L \)-positions can be described roughly via \( N = N_0 e^{-E/k_B T} \) – a feature well known from electron beams generated by LPAs. In LPA research, \( k_B T = T_{\text{eff}} \) is called the effective
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Figure 39: Van Allen belt spectral electron flux according to NASA’s AE8_{min} model [4] at distances from Earth of \( L = 1 - 10 \). In a 2D visualization in a), the orbits of GPS and Galileo satellites as well as the geostationary (GEO) orbit are indicated, while the 3D visualization b) indicates the exponential distribution of the spectral flux (from [5]).

temperature of the beam. This indicates that the directed electron flux on stationary orbits can be reproduced by exponential electron beams from LPAs, by varying the temperature of the LPA-generated electron beam. This would constitute a fundamental advance, since the current state-of-the-art in radiation testing does so far not involve reproducing the exponential electron flux – simply because particles sources such as LPAs which enable controlled, tunable exponential-energy flux have hitherto not been known in the radiation testing community – but instead uses more simplified methods, making use of more conventional particle sources. For example, one basic and widely used approach is to use radiation generated by radioactive sources such as \(^{60}\)Co (\( \gamma \)-rays and electrons) to evaluate total-dose effects. By varying the distance from the radioactive source and/or the time of irradiation, the total dose received by the device under test (DUT) is tuned. In addition, to produce electron flux of a certain monoenergetic energy, linacs are widely used. However, such monoenergetic electron flux does not exist in space. For this application, the ability of LPAs to produce flux with exponential energy distribution such as present in the radiation belts is therefore highly desirable.

Figure 40 shows explicitly, that at certain fixed orbits, the radiation belt electron flux follows an exponential decrease. As an example, the electron flux at \( L = 3.17 \) (GPS), \( L = 3.65 \) (Galileo), and \( L = 6.65 \) (GEO) are plotted. By fitting an exponential decay function, effective temperatures \( T_{\text{eff}} \approx 0.4 - 0.62 \) MeV can be deduced.

Monoenergetic beams are not well-suited to reproduce the specific features of such
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Figure 40: Electron flux according to AE8_{max} at the important orbital distances of GPS, Galileo (both around medium Earth orbits (MEO)) and at geosynchronous orbit (GEO). Fitting exponential particle distributions leads to temperatures in the range of \( T_{\text{eff}} \approx 0.4 - 0.62 \text{ MeV} \) (from [5]).

Flux with roughly exponential energy distribution. This is true for flux incident on the space vessel, and also after passing through thin layers of matter – which corresponds to spacecraft shielding and therefore is a very relevant scenario. To illustrate and underline this, figure 41 a) shows how a monoenergetic beam of energy \( E = 4 \text{ MeV} \) is decelerated and damped when passing through matter. The flux has been calculated making use of the Monte Carlo code MULASSIS [259], which was developed explicitly for radiation shielding analysis. The marks {\text{i), ii), iii) and iv)) correspond to the flux after 1, 2, 3, and 4 mm of aluminum. In contrast, 41 b) shows how the flux changes when an exponential-energy beam with \( T_{\text{eff}} = 2 \text{ MeV} \) is incident on the same shielding. The spectral flux is damped, but apart from the low-energy cutoff, still has an exponential shape. This behavior is complemented by the results of total ionizing dose (TID) calculations which have been carried through for the monoenergetic and the exponential-energy beam, and are given in figure 41 c) and d). The TID increases during passage through the solid aluminum in case of the 4 MeV beam (c), and reaches its maximum after about 3 mm, when a large fraction of particles has been decelerated down to low energies, and their energy is deposited. In contrast, with the exponential energy flux (d) the TID decreases constantly during passage through matter, which is the realistic case expected for radiation belt electrons, too. These very fundamental considerations show that exponential-energy beams are much better suited when compared to monoenergetic beams to reproduce radiation belt electron flux, and thus would enable to develop advanced radiation hardness testing procedures and standards. LPAs are ideally suited as
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Figure 41: Comparison of spectral flux change of a) an incident monoenergetic electron beam (as from a linac) of $E = 4$ MeV when straggling through matter and b) an exponential beam as from an LPA with $T_{\text{eff}} = 2$ MeV. The marks i) to iv) denote the forward flux after passing through 1 to 4 mm of aluminum. In contrast to the monoenergetic beam, the exponential-energy beam remains exponential with the exception of the low energy cut off (from [5]).

such radiation sources.

Electrons in the outer van Allen belt have maximum energies of up to $E \approx 10$ MeV. Laser interaction with overdense targets can produce reliably exponential energy electrons in this energy regime with high charge, and is therefore highly suited for this scenario. Furthermore, scaling laws are well-known which predict the electron temperature as a function of laser intensity $I$. These scalings are subject of research since decades and describe the dependence of non-relativistic [85, 98, 99] as well as of relativistic laser intensities [100]. Empirically, a power-law scaling $T_{\text{eff}} \propto (I\lambda^2)^\zeta$, where $\lambda$ is the laser wavelength, describes the effective electron temperature where $\zeta$ typically ranges between 1/2 and 1/3 [89]. According to [101], $\zeta = 1/2$ in the intensity range of $I\lambda^2 \approx 1.3 \times 10^{18}$ W/µm²/cm² to $I\lambda^2 \approx 1.4 \times 10^{19}$ W/µm²/cm², resulting in the explicit scaling which can be expressed as $1.37 \times 10^{18}$ W/cm²/λ²$(T_{\text{eff}}[\text{MeV}]/m_0c^4 + 2T_{\text{eff}}[\text{MeV}]/m_0c^2) = I[\text{W/cm}^2]$ and defines an intensity working point suitable to pro-
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produce flux with the aim temperature. In another well-known work in the intensity range from some $10^{16}$ W/cm$^2$ up to $10^{19}$ W/cm$^2$ [102], a value of $\zeta = 1/3$ is inferred, leading to a scaling $T_{\text{eff}}[\text{MeV}]^3/\lambda^2 = I \times 10^{20}$ W$\mu$m$^2$/cm$^2$.

Now, to produce electron flux with an electron temperature of $T_{\text{eff}} \approx 0.35$ MeV, according to radiation belt flux at $L = 3.5$, the laser intensity on target (assuming that a Ti:Sa laser pulse with a central wavelength of about $\lambda = 0.8 \mu$m is used) should amount to values between $I_{L=3.5} \approx 3.9 \times 10^{18}$ W/cm$^2$ ($\zeta = 1/2$) and $I_{L=3.5} \approx 6.7 \times 10^{18}$ W/cm$^2$ ($\zeta = 1/3$).

By integrating the differential flux according to NASA AE8, the total fluence on a vessel can be calculated and is the reference for LPA-based testing campaigns. For example, a total number of $\approx 3 \times 10^{12}$ cm$^{-2}$ electrons can be estimated to be incident per day on a vessel operating at $L = 3.5$ during solar maximum activity. A corresponding practical scenario would be a satellite close to GPS orbits, where the electron flux can increase by about one or two orders of magnitude in rise times of the order of 1-2 days, for example [260]. A typical total charge of $Q \approx 100$ nC can be assumed to be emitted in the forward direction per shot by LPA interaction, corresponding to 3.5% of the energy of a 1-J laser pulse, which is a typical energy value for the LPA driver pulse. At an intensity of about $I_{L=3.5} \approx 5 \times 10^{18}$ W$\mu$m$^2$/cm$^2$, producing electron flux with a temperature of $T_{\text{eff}} \approx 0.35$ MeV one would need only $\approx 5$ laser shots, or half a second of LPA performance at 10 Hz repetition rate, in order to produce the total flux incident on a vessel at $L = 3.5$ in space per day. However, this value is only theoretical and underestimates the total number of shots needed, since one cannot put the DUT directly behind the target, and because the electron flux generated via laser-solid interaction is emitted in a cone with broad divergence. The divergence is also intensity-dependent [13, 104], and amounts for the above estimated intensity range to some 25$^\circ$ full divergence angle. Assuming a radially Gaussian intensity distribution, a DUT with an area of 1 cm$^2$, positioned at a distance of 10 cm away from the radiation source on axis, would need 56 shots in order to receive the maximum daily dosis of a satellite surface. Similar to what is done in conventional radiation hardness testing with $^{60}$Co sources, the received flux can be tuned by varying the distance.

As a further remark, laser systems are currently under development which promise to produce multi-mJ energies at kHz repetition rates. Such systems would be focussable to intensities in the $10^{18}$ W/cm$^2$ regime. With such systems, the yearly flux on a satellite could be reproduced in the laboratory within seconds. Similarly, it is conceivable to use a high-power, 10-Hz system, and to split the pulse into several delayed laser beams in order to increase the effective repetition rate.

There is a fundamental difference between radiation belt flux and LPA-generated flux. Radiation belt flux is quasi-continuous, whereas LPA-generated electron flux is initially pulsed, since it is generated during the laser pulse interaction with plasma electrons. At the source, the duration of the individual electron beams is equal to the laser pulse duration. However, since we have exponential energy distribution, the time of flight
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Figure 42: Reduction of exponential-energy electron flux due to energy-dependent velocities and divergence. In a), the flux of a beam with $T_{\text{eff}} = 0.35$ MeV, $Q = 100$ nC and a divergence of $\theta = 25^\circ$ through a DUT area of 1 cm$^2$ is calculated at distances 0.1 cm, 1 cm, and 10 cm behind target (note the logarithmic scaling). Next, the influence of the divergence is visualized by plotting the flux through 1 cm$^2$ after a distance 1 cm (b) and 10 cm (c) for the beam with parameters as in a), but for a hypothetical divergence of $\theta = 0^\circ$ and $\theta = 25^\circ$ (from [5]).

The pulsed nature of LPA flux in combination with the tunable peak flux due to divergence and exponential energy distribution has positive as well as negative aspects. On the one hand, too high peak flux might lead to "unnatural" effects which would not occur in the quasi-continuous, comparably low-flux environment in the radiation belts. On the other hand, the ability to increase flux to high peak levels enables to study the threshold at which such nonlinear effects might occur. This can be also advantageous in the context of both natural as well as unnatural rapid flux release events.
9.1 Experimental setup for Earth radiation belt reproduction and electronics testing

For our proof-of-concept experiments, we used the Arcturus laser system at Heinrich-Heine-University Düsseldorf, and the proton cyclotron at UCL for calibration of certain image plate diagnostics.

The Arcturus laser system at the Institut fuer Laser- und Plasmaphysik of the Heinrich-Heine University is a commercial 150+ TW chirped pulse amplification system from Amplitude Technologies. The gain medium is a Ti:Sapphire crystal (Ti$^{3+}$:Al$_2$O$_3$). The Ti:Sapphire oscillator provides pulse of 23 fs duration, with an energy of barely 5 nJ and a repetition rate of 76 MHz. To provide such short pulses the bandwidth is 90 nm at a central wavelength of 800 nm. The initial repetition rate is reduced to 10 Hz by a Pockels cell. Before the pulse amplification the pulse is temporally stretched to 300 ps by a stretch grating.

The amplification of the pulse occurs in several steps. First the pulse is coupled into a resonator of a regenerative amplifier, where it gains energy up to 0.5 mJ. In addition to the gain the laser pulse mode is set here. As in all following amplifiers a frequency double Nd:YAG laser with its central wavelength at 532 nm is used to pump the Ti:Sapphire amplifier crystal. In the further process the pulse gains energy at in total three multi-pass
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amplifiers and reaches an energy of 2.7 J at maximum just before it enters the compressor section of the system. To not exceed the damage threshold of the optical components, the diameter of the beam is increased after the passage of the first amplifier to 8 cm. The main amplifier crystal is cooled in a cryogenic section to suppress the formation of a thermal lens. After the amplification process the chirped pulse is temporally compressed back to 23 fs by the a grating compressor, to compensate the frequency depending run time. The intensity of the pulse \( I = 10^{12} \text{W/cm}^2 \) is already enough to cause nonlinear processes in air such as ionization and the optical Kerr-effect. This would lead to a degradation of the spatial and temporal beam profile. Hence the pulse compression and further propagation is performed under high vacuum conditions \( p < 10^{-5} \text{mbar} \). Because of energy loss at the optical elements in the compressor and on the pathway to the target (absorption, scattering, grating efficiency) roughly 40% of the amplified energy, here 1.2 J, are deployed at the target. Figure 43 shows a schematic of the laser system. Details of the system after comissioning are given in [261].

To achieve the required intensity levels, focusing of the laser beam is done with a F/2 off-axis parabola with a focal length of 15 cm. This results in a Rayleigh length of 30 \( \mu \text{m} \). The hard focussing enables a spot size of a few square micrometer and therefore intensities in the range of \( 10^{19} \) to \( 10^{20} \text{W/cm}^2 \) for the given pulse energy and duration. Figure 44 shows a typical focus spot as observed with the focus diagnostics during the experimental campaigns. Such pictures are taken at reduced laser intensity (ASE level) and with a lot of filters, because otherwise the CCD chip would be overexposed.

**Figure 44**: Focal laser spot as recorded at reduced intensity with a microscope objective and a CCD camera.

To produce the required van-Allen belt class electron flux, we chose TNSA-type [132] laser-solid interaction in the intensity range of about \( I \approx 5 \times 10^{19} \text{W/cm}^2 \), which is a typical intensity value for state-of-the-art laser-plasma-accelerators and can easily be reached by commercially available systems such as the Arcturus laser system. The target material was a 30.6 \( \mu \text{m} \) thick aluminum foil. The thickness was chosen to enable a high number of laser pulses without changing the target foil in this scenario. Much
thinner foils tend to tear. The holder is a frame structure with one open side, to have a free field of view onto the target parallel. The target can be moved into all spatial directions by three linear positioning systems with 1\( \mu m \) accuracy. The laser incidence is 45\( ^\circ \). To change the angle of incidence that target holder is attached onto a one-axis mount, which allowed a rotation of \( \pm 30^\circ \).

While the interaction of the laser pulse with the target foil produces particle beams within the duration of the laser pulse on the femtosecond scale, the energy deposited by the laser pulse leads to melting of the material after each shot. The resulting holes in the target foils are shown in figure 45.

Figure 45: Holes in the target foil – each hole corresponds to one laser shot and is generated by melting due to the deposited laser energy after the laser-plasma interaction.

Figure 46: The target foil is getting bumpy when one comes close to the edge of the target foil. This should be avoided because the bumpyness is of the order or larger than the Rayleigh length especially with particularly thin foils.

Figure 47 shows the complete setup inside the target chamber. The compressed laser pulse is incident from the right hand side and then is send to the 90\(^\circ\) focusing parabola.
Figure 47: Setup inside the irradiation chamber. The incident laser system is strongly focused on an Aluminum foil target, where the radiation is produced. Focus diagnostic microscope objective, image plate (IP) stack and permanent magnet based spectrometers in forward and backward direction are shown next to the target foil positioning system.

The F/2 parabola then focuses the laser beam on the target foil, which was moved after each shot to provide a fresh surface for the interaction process. Before the laser pulse reaches the target foil, no radiation is produced. An image plate (IP) stack, which also can hold the devices under test (DUT’s), to record the electron flux in forward direction is put on axis. On axis, the stack has a central hole where electrons (and protons) could pass and enter the magnetic spectrometer for further energy measurement.

FUJI image plates are extensively used as diagnostics. Originally being developed primarily for x-ray detection (e.g. for medical use), they are a well-established diagnostics tool in the laser-plasma-accelerator world. They have been cross-calibrated and absolutely calibrated in earlier works [8], and have proven to be very reliable tools in
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Figure 48: Image plates (FUJI). Incident electrons (or protons/ions/photons, depending on their energy) can excite metastable states during irradiation which can then be detected in a scanner system, where a HeNe laser beam scans the image plate. The red laser light relaxes the metastable states, and the emitted radiation is collected via photomultipliers to generate the data.

harsh laser-plasma-environments. IPs have an extremely high linear dynamic range of $10^5$ and can be scanned with high resolution and a pixel size as small as 25 $\mu$m. There are various types of FUJI image plates available, which however share the general concept and composition. This is shown in figure 48. A protective layer (a) of Mylar is covering the sensitive phosphor layer (b) consisting of active BaFBr:Eu$^{2+}$ molecule centers in a stable emulsion layer, followed by an intermediate layer (c) and a base layer (d) and a weakly magnetic layer (e) which is used to fix the flexible plate during the scanning process.

After irradiation, the signal is stored in the excited metastable states until these states are relaxed/erased by light or over time. In the scanner, a well-steered relaxation process takes place in which a HeNe laser ($\lambda = 633$ nm) beam scans the plate line by line. When the laser beam hits excited metastable states, these states are relaxed, and higher wavelength light ($\approx 390$ nm) is emitted, which is collected in a photomultiplier and is then attributed to the position of the scanning laser beam on the IP. The sensitivity is high enough to detect the signatures of single electrons.

The dynamic range is larger than the intensity range which can be read out in a single scan process with the scanner system (Fuji BAS-1800 II). At high signal intensities it is necessary to perform several scan processes consecutively. The decrease of the signal intensity depends only on the number of successive scans and not on the initial intensity. The erase rate was determined by calibration with an $\alpha$-emitter ($^{226}$Ra, $A = 3.3 \cdot 10^5$ Bq,
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$T_{1/2} = 1602 \text{ years}$) in the context of the present R&D activity. The exposition times were varied between 10, 100, 1000, and 3600 s. Every image plate was read out consecutively until no "overexposed" areas were measured anymore. The initial signal intensity can be reverse determined by the number of scans. Figure 49 shows the averaged measured erase rates $L_n$ per scan $n$, which can be well approximated by the exponential fit

$$L_n^{2012} = 4.27 + 25.59 \times \exp \left(-\frac{n}{4.99}\right).$$ (113)

Figure 49: Averaged erase rate of the image plate signal as function of number of scans and the related exponential fit.

To determine the number of electrons from the measured photo-stimulated luminescence (PSL) values, the calibration by Tanaka [7] and the extension by Hidding [8, 9] is used. Figure 50 illustrates the calibration curve of PSL per electron as a function of the electron energy for different types of image plates. For the proof-of-concept experiments mostly the BAS-TR type is used. The ratio of the measured PSL signal and PSL per electron gives the electron number.

More details on image plates and their calibration are given in [8]. These image plates are used as a detection screen in permanent magnet spectrometers, for cross-calibration with lanex screens, and in a stack method first described in [8].

During the laser-plasma-interaction, various processes lead to emission of broadband electron radiation in various directions. This is highly advantageous, because this allows in principle for harvesting this broadband particle flux for irradiation of additional test devices. It also improves very much the efficiency of the whole process as regards wall
Figure 50: IP signal dependency of the electron energy. For different IP types Monte-Carlo simulations (solid line, left y-axis) and the PSL calibration curve in accordance with Tanaka [7] is shown. This is based on three measurements obtained in conventional accelerators for electron energies of 11.5, 30, and 100 MeV and earlier measurement results for the lower energy range. Taken from [8,9].
plug power to usable radiation for testing. To explore these options, the experimental setup involved diagnostics, as well as irradiation of test devices, in a large range of directions, such as in the target normal direction behind the target (dubbed target normal forward direction), in the laser beam propagation direction (dubbed laser forward direction). The particle emission direction in laser forward direction differs from the target normal forward direction because the laser incidence angle on target can and should never be perpendicular, because then there is the danger of backscattering of a significant amount of laser energy back into the beamline and potentially, and even more dangerous, into the laser system itself where it could damage optical elements as well as the laser amplifier medium itself. Especially when the laser incidence angle is low (measured in respect to the target normal direction as usually in optics), the emission cone resulting from acceleration processes in the laser forward direction and the emission cone resulting from the acceleration process in the target normal forward direction do overlap, which can make the monitoring and interpretation of observed emission complex. In addition to these two forward emission angles, there are other processes which lead to acceleration of electrons in the target foil parallel directions, dubbed target parallel forward direction and target parallel backward direction, respectively. The peculiar differences and similarities in these acceleration processes are described in 7 in detail. Figure 51 gives an overview on the different radiation emission directions in polar coordinates, the used nomenclature and the positions of the main diagnostics. This schematic figure relates to figure 47, the photo of the actual experimental setup inside the chamber.

The five different main emission directions with a laser incidence on target of 45° are:

180°: laser forward direction
135°: target normal forward direction
315°: target normal backward direction
225°: target parallel forward direction
45°: target parallel backward direction

Highlighted here are the incident laser beam (a), the particle emission cone in the target forward parallel direction (b) colored in blue, the target normal forward direction particle emission cone (c), also colored in blue, the target itself, here an aluminum foil in a holder sized 4 cm × 4 cm × 30.6 µm (d), the combined DUT holder and image plate stack (e), a permanent magnet based spectrometer which deflects electrons according to their energy in the upwards direction, where they are then detected either by an image plate, a fluorescent screen for online diagnostics, or a combination of both (f), and an image plate stack to diagnose the target normal parallel forward electrons (g).

For characterization and monitoring of the emitted electron and proton flux in the various directions a wide range of diagnostics was used. The imaging plate based sandwich
Figure 51: Conventions and naming scheme for the detection directions and position of main diagnostics in the target chamber.
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Figure 52: Sketch of setup at low incidence angle. The laser (red) is incident on target and accelerates electrons in target normal forward direction and in parallel direction (blue).

Stacks developed in [8] proved very useful for large area monitoring and for radiography images during the irradiation of electronic test devices. We used large imaging plate stacks with an area of up to 15x15 cm and up to 7 image plates in one stack, and several miniature image plate stacks which were developed in the context of the project. An addition to the imaging plate stacks, up to three dipole magnet spectrometers with different sizes are used. While image plates are great tools for recording electron (and proton) flux over many or single shots, one fundamental disadvantage is that they have to be taken out of the vacuum chamber before the readout process. Therefore one the one hand, as a complementary online diagnostics phosphorizing beam viewers which were observed with triggered cameras were used. These beam viewers were used in combination with imaging plates (for cross-calibration) as well as as standalone diagnostics to monitor the spatially resolved electron flux per shot. On the other hand, a novel technique of vacuum insertion boxes was developed. Extension boxes are quite commonly used in laser-plasma-experiments to expand the vacuum chamber locally, for example in case a specially bulky diagnostics is needed. In turn, our goal was a) to get closer to the radiation source (i.e. the target foil) and b) to have locally a much thinner wall when compared to the vacuum chamber itself, so that the particle flux is not completely stopped. This was achieved by using a 2 mm thick Al sheet to separate the round insertion box from the vacuum. Figure 53 shows the insertion box as integrated into the vacuum chamber. On the left hand side of figure 53, one can see the cylindrically shaped insertion box mounted inside a rectangular chamber wall. The end face of the box is a 1 mm thin Aluminum sheet (in contrast the standard vacuum...
chamber walls have a thickness of \( \approx 3 \) cm, which at these energy levels nearly completely blocks the particle radiation) with an areal diameter of approximately 30 cm. There is a miniature image plate stack mounted to this face, which consists of alternating layers of small image plates and 5 ct coins. A zoom on such a stack is shown as an inset in the figure the to left, where one can see the coins and a round piece of image plate on top. Because of the light sensitivity of the stacks, all the image plate stacks are covered with aluminum foil (as used for domestic needs), which is seen in the figure on the left and in the middle, where it is zoomed on the insertion box end face. DUT’s can be mounted either on the vacuum side of this end face or on the outer face. In order to increase the flux on the DUT’s, it will be in most cases desirable to mount the DUT’s in the inside of the vacuum chamber, but in case of larger DUT’s and/or actively powered electronics such insertion boxes offer an elegant way to get close to the source, where the flux and fluence are higher, without the need to work in vacuum. In addition to the mini image plate stacks, we have in addition partially used large image plate stacks on the outside of the insertion box face. Such large stacks allow to measure the flux over the whole insertion box face, and can be removed for readout either after a single or a few shots.

The use of the insertion box and diagnostics such as the image plate stacks outside the vacuum chamber is a considerable advantage when using image plates only inside the vacuum because then the vacuum does not have to be broken, which results, depending on pumping capabilities, in an interruption of up to an hour betweens irradiation cycles. However, even with image plates outside the bunker the operation procedures require significant time for entering the radiation bunker, for removal of irradiated IPs and for the placement of new IPs, and for the readout process, saving of the data and erasing of the residual on IP. While the latter can of course be done later or during the next irradiation cycle, the required time still amounts for a couple of minutes. It is therefore
used mostly only as sample diagnostics, e.g. when important interaction parameters are changed and especially thorough control diagnostics is desirable.

In addition to the high-resolution long-term diagnostics such as integrating image plate stacks behind the DUT’s inside the vacuum chamber (compare figure 60), medium-term diagnostics such as image plate stacks outside the chamber in the insertion box, we have also employed online diagnostics which allow for shot-to-shot measurement. These are based on fluorescent screens instead of image plates. Such (lanex) screens are often used in the laser-plasma-accelerator community as well as in the conventional accelerator community. The incident electrons generate the fluorescence signal, which can then be monitored with triggered cameras observing the screens. The screen size is comparable to those of image plates, and can therefore cover large areas. We have extensively tested various kinds of fluorescent screens in a previous project, see [262] for details.

In our proof-of-concept campaigns, such fluorescent screens are used both as alternatives and add-ons to image plates, both in the detection plane of magnet spectrometers as well as on the back side of image plates such as those in the insertion box. For example, the CCD camera (red) figure 53 on the left is used to monitor a fluorescent screen which is placed behind the mini IP stack, DUT’s and image plate stacks (all optional), also see figure 60, right hand side.

In addition to the main emission directions, the position of the main diagnostics is also indicated in the schematic figure 51. To characterize the emitted electrons and to get information on the applied dosage in the DUT’s (here: optocouplers) IP stacks and behind these spectrometers (large: target normal forward direction, smaller: target normal backward direction) are placed, while in target parallel directions only image plate stacks and mini stacks are placed, and in laser forward direction the geometry of the setup allowed mostly only for use of mini IP stacks. In addition to what has been said above, a very small permanent magnet spectrometer could be placed inside the insertion box, to allow for online diagnostics of the electron beam energy distribution emitted in this direction.

To go into more detail, figure 54 and 55 show the setup for testing and monitoring the electron flux as used in the target normal forward direction. In figure 54a), a compact view is presented, which shows the IP stack (which optionally can contain DUT’s in the front), and the appended permanent magnet based spectrometer behind it. In the detection plane, either one (ore more) IPs are placed, or a fluorescent screen, or both. Figure 54b) shows an exploded view of the first part of the setup, the IP stack. A first layer of variable thickness acts as proton shield, after that follows the DUT section. In the case of simple components the DUT’s can be devices of the order or 1 cm$^2$ or even smaller. In this case, the DUT’s do not block completely the large-divergence electron flux (indicated by the green cone), so that it is possible to monitor the electron flux behind the DUT layer with an electron-sensitive image plate (IP) stack similar as in [110] or in [263].

Such a stack consists of image plates with a well-defined response function for elec-
Figure 54: Experimental setup to detect electron flux and to place optocoupler devices. 

a) The exponential electron and proton flux from the target is incident on an IP stack with or without DUT’s, and the on axis fraction of the electrons then enters the magnetic spectrometer through a small hole and is resolved in energy. b) Exploded view of the IP stack box, revealing the sandwich composition of IP/stopping layers (from [5]).
Figure 55: Variation of the setup of figure 54, where the front more bulky shield is replaced by a frame on which foils of various thickness could be placed.

Figure 56: Photos of the permanent magnet spectrometer, corresponding to the analysis and calibration in figure 57.
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Figure 57: Simulation results for electrons up to 30 MeV. The particle trajectories and the dispersion curves are shown.

electrons [7,264] alternating with beam attenuator and stopping layers, for example, simple aluminum plates. As discussed above in detail, the image plates are reusable and have a large dynamic range ($> 10^5$) to be used to accumulate many shots without being overexposed, so that they can be scanned after the irradiation to read out the accumulated electron signal. Monte-Carlo type simulations similar to those used to produce figure 6 were then used to reconstruct the incident electron beam flux and temperature. It is especially helpful that the electron sensitivity response of the IPs is nearly constant for electron energies $> 1$ MeV (compare figure 50), because the proton shielding does also block $< 1$ MeV electrons and facilitates the reconstruction of the signal incident on the DUT's.

In practice it has been found that a stack consisting five image plates is sufficient for the expected electron energy in the used intensity range. The low energy electrons will be absorbed in the front IPs, the high-energy electron in the rear (compare figure 69. The aluminum attenuators between the IPs are 1 mm in thickness. In normal forward direction the attenuators and IPs have a size of $10 \times 10$ cm, and $12 \times 12$ cm in normal backward and parallel backward direction, respectively. The complete stack as well as the DUT holding plate have a small hole on axis, as is indicated in figure 54b). With the exception of the front proton shielding layer, the incident on-axis fraction of the electron beam can pass the setup without being attenuated.

Behind the image plate stacks the electrons enter a dipole magnet spectrometer.
proof-of-concept experiments

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Figure 58: Example of raw data of spectra as recorded on the image plates (left) and lanex (right) in the detection plane of the permanent magnet based spectrometer in forward direction.

through a mm-scale aperture to be further analyzed. Electrons are dispersed according to their energy in a permanent magnetic field and are detected by a combination of IP and scintillating screen in the detection plane. The remaining part of the electron beam is dumped in a low-Z-high-Z sandwich dump, so that these electrons do not falsify the signal as recorded on the energy detection plane IP. This IP can after irradiation be scanned to give further information on the average energy distribution of shots. In addition, a calibrated scintillating screen [265] directly behind this IP is used to monitor the energy distribution and intensity and potential shot-to-shot variation online with a triggered CCD camera.

The spectrometers are based on hard ferrite magnets consisting of barium and strontium ferrite (BaFe, SrFe) with a magnetic remanence of \( B_r = 380 - 400 \, \text{mT} \). Both spectrometers use an U-shaped yoke of magnetized steel (ST-37) to generate a horizontal and broadly homogeneous magnetic field. The obtained magnetic flux density is 0.23 T with the small spectrometer and 0.29 T with the largest one, respectively. Figure 56 shows photos of the smaller spectrometer (the larger one has an analogous composition), showing the magnets, the yoke, the front blocker section, and a lab jack for in-vacuum use. Electrons accelerated in both target normal directions are passing on the beam axis the aperture of the spectrometers and are directed due to the Lorentz force in the detection plane, which is parallel to the incident electron axis. The energy distribution is detected by an Imaging Plate. The Imaging Plate is covered by a fluorescent screen to enable online shot-to-shot measurements. The front side of each spectrometer is covered by a combination of a low-Z-material (2.5 cm Polyacetal (POM)) and a high-Z-material (1.5 cm lead) to avoid scattering of electrons not passing the aper-
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ture. The influence of plasma glowing and scattered laser light was minimized by black aluminum foil wrapped around the spectrometer sides. To determine the nonlinear dispersion both spectrometers were modeled in CST PARTICLE STUDIO and the particle trajectories were simulated. Fig. 57 gives the deviation in the plane of detection for both spectrometers. The small spectrometer was used in normal backward direction, the large spectrometer in normal forward direction.

In figure 59 it is shown schematically how the raw spectra displayed in figure 58 are recorded on the Lanex beam viewer. The CCD camera looks from outside the vacuum chamber through a window on the Lanex, using a flat mirror inside the vacuum chamber to reflect the fluorescent signal, which is generated by the electrons of different energies arriving at different positions on the Lanex after deflection in the magnetic field of the spectrometer, and after passing the IP layer.

Figure 59: Schematic view of online measurement of the electron spectral flux on axis using a lanex screen observed via a mirror and a triggerable CCD from outside the vacuum chamber.

Next, figure 60 shows how DUT’s (here: high-quality optical dielectric mirrors for use in space in the context of gravitational aves detection) are mounted to the inner wall of the insertion box inside the vacuum box, using tape band. Indicated with the greenish ellipses is the main flux of electrons in target parallel direction due to the TSPA mechanism. One of the ellipses stems from emission into the target front side acceleration processes, the other from acceleration on the target rear side. In between the both main emission directions is a gap due to the exact target parallel direction. The target foil is and the target holder are not seen in the picture because they have been removed to provide clear view on the insertion box face. In the actual irradiation experiments, image plates and fluorescent screens were placed on the other side of the insertion box wall, compare figure 53. Depicted on the right hand side of figure 60 are two example shots of such images of the electron flux on the lanex screen outside of the chamber.
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The elliptic structure of the electron flux (mirrored because the CCD is looking from the other side) is again seen clearly. In the shot in the bottom of the right hand side one sees in blue on the left hand side the round shadow of an image plate stack (compare again figure 53).

Figure 60: Left: DUT’s (optocouplers and mirrors) mounted to the the inside wall of a vacuum insertion box. Right: Signal data on beam viewer observed with a CCD cam on the other side of insertion box (at air) after irradiation.

The combination of the described diagnostics enables to have full information on the incident electron flux. The effect of tuning the electron temperature via changing the laser intensity, if necessary, can be directly observed. Since the electron flux can be assumed to be radially symmetric, the DUTs should be placed close to the axis on concentric rings, with enough space between the single DUTs to allow for total flux reconstruction after the irradiation by analyzing the IP stack. Finally, it should be mentioned that the temperature slightly off axis will be slightly lower than on axis, since
the most energetic electrons are emitted in the forward direction on axis. The off-axis temperature dependency can be checked by rotating the whole setup around the electron beam axis.

9.2 Measured particle flux data and characterization of effects on DUT’s

In the first phase of the experimental campaigns we carried through extensive calibration runs to test the diagnostics and to find suitable parameter regimes for our goal of reproducing van-Allen belt space radiation for testing purposes for the first time in the laboratory here on Earth. After these runs, we inserted in agreement with European Space Agency optocouplers from various manufacturers in front of the image plate stack, behind the proton shielding foil. Photos of these optocouplers are found in figure 61.

Optocouplers as DUT’s fulfil three main requirements we had for the proof-of-concept experiments. They should a) be relevant for space electronics, b) be able to be tested in a passive setup (without the need to use them onboard of a powered electronic circuit), be rather compact. We have used different kinds of optocouplers in close consultation with ESA:

Vishay SFH6345 (8 pins)

![Vishay SFH6345 - 8 Pins](image1)

Fairchild CNY17F-3 - 6 Pins

![Fairchild CNY17F-3 - 6 Pins](image2)

Liteon 355T - 4 Pin Gull Wing

![Liteon 355T - 4 Pin Gull Wing](image3)

Isolink OLH400/300/249

6 Pins mit TO-5 Fassung

![Isolink OLH400/300/249](image4)

Isolink OLS449

auf DIL-Fassung

![Isolink OLS449](image5)

Micropac 6626-101

![Micropac 6626-101](image6)

Figure 61: Optocouplers from various manufacturers and different makes used in the proof-of-concept irradiation campaigns.
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Fairchild CNY17F-3 (6 pins)
Liteon 355T (4 pins gull wing)
Isolink OLH400/300/249 (6 pins w/ TO-5 base)
Isolink OLS449 (8 pins)
Micropac 6626-101 (8 pins)

Various means of holding mechanisms were discussed for these and other devices, such as fixing them with screws on a small breadboard (see figure 62), but in the end a more flexible yet simple method was chosen. The devices were simply clipped on anodized thin (≈ 50 microns) aluminum foil, as shown in figure 63. This has a number of advantages: first, the optocoupler pins do not have to be bended in order to fix them to the carrier plate, and second, there are no screws which would otherwise throw rather bulky shadows on the radiography images (see next sections). The aluminum foil is thick enough to safely hold the optocouplers, and flexible and thin enough to stick the optocoupler pins through, and to add no significant radiation stopping before the diagnostics image plates are reached.

Figure 62: Mounting of optocuplers on a miniature breadboard. Not chosen in current experiments but a potentially viable option for other DUTs.

Figure 64 visualizes how the optocouplers in the Al foil holder are then inserted into the IP stack box (left hand side of figure 64) and in front of a simpler, more compact image plate stack (right hand side of figure 64). In any case, the optocoupler layer is covered by at least another layer of anodized Aluminum to protect them from low energy protons, potential debris and reflected laser light.

In figure 65, the temperatures obtained in the target normal forward and backward directions in the magnetic spectrometers. Compared are here the obtained spectra data as recorded on the IPs in the spectrometer detection plane in target normal forward and backward (black solid square/green solid circle, respectively) directions, data observed
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Figure 63: Elegant mounting of optocouplers by spiking them on a thin anodized aluminum foil. Radiation stopping is therefore minimized and a clear picture is seen on the following image plates and other diagnostics. The central hole is for particle detection on axis.

Figure 64: Optocouplers mounted al Al foil and inserted in IP stack box (left) and in front of a more simple stack to harvest another radiation emission direction (right). In any case the DUTs are covered by an additional anodized Al foil or even thicker proton shielding.
online on the lanex screens in the same plane (black hollow square/green hollow circle, respectively). The laser intensity is scanned, and generally speaking, the obtained electron temperature is thus increased as laser intensity grows. This holds both for forward as well as for backward directions. For most data points, it is averaged over many tens or even hundreds of shots, while at an intensity of $I \approx 6.5 \times 10^{19}$ W/cm$^2$ only one single shot has been recorded. Many more shots have been taken with the lanex screen alone (without an IP inserted in the detection plane beneath the lanex screen, see figure 59.

![Figure 65: Comparison of temperatures as obtained by the IPs and Lanex detectors, showing good agreement.](image)

Figure 66 and table 2 show a direct detailed comparison between a spectrum obtained at $I \approx 4.5$ W/cm$^2$ with an IP (red line) and the lanex viewer and CCD cam (black line). The resulting retrieved temperatures shows very good agreement with both methods. Note that the left $y$-axis is for the IP, while the right $y$-axis is for the Lanex – and both have very different scaling, so while at first sight the agreement seems to be not good, it actually is, as can be seen from the retrieved temperature values, namely $k_B T = T_{\text{eff}} \approx 0.54$ MeV in case of the IP, and $T_{\text{eff}} \approx 0.57$ MeV in case of the Lanex (again, the temperature can be used via $N = N_0 e^{-E/k_B T}$ to characterize the spectrum). As regards noted differences, the lanex signal is rather noisy at electron energies above approximately 2 MeV – this is attributed to the lower sensitivity of the lanex on the one hand and the limited quality of the CCD camera (8 bit). This shows vividly that image plates are better suited when one needs to retrieve highly accurate data, while Lanex screens have the big advantage of being online-ready diagnostics. It shall be emphasized here that the use of image plates is a relatively novel method which is just in the process of being introduced in the conventional accelerator field – one of the increasing number of cases where a diagnostics has been first introduced in the laser-plasma accelerator community and then is transferred to the conventional (rf-cavity based) accelerator field, and not the other way around as is the case for most diagnostics. Since fluorescent beam
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The total electron yield per shot per steradian is an important number. Image plates (or image plate stacks) are ideal means to accurately record the electron flux over many shots. These image plates were placed at a distance of 5 cm from the target, such that one square centimeter corresponds to 0.04 sr. In Figure 65, electron numbers are given per square centimeter. The resulting conversion factor from square centimeter is 25. The total electron yield was measured in various directions. Shown in figure 67 are intensity scans of the electron number per shot per square centimeter. On the left hand side of figure 67, the electron yield in target normal forward direction is shown. There is a clear minimum of the electron yield at an intensity of about $I \approx 4.7 \times 10^{19} \text{ W/cm}^2$. At larger intensities, there is a trend of increasing electron yield. Physically, this trend is partially attributed to the different laser absorption and acceleration mechanisms which change as the intensity changes, refer to section 7. Also, with a protecting aluminum foil thickness of 200 µm as used in this case in the forward direction, one is not yet beyond the peak in the sensitivity response of the image plates, compare 50, which complicates the accurate attribution of electron flux numbers. The electron number in case of the single shot (red) is much higher than in the averaged shots, which makes it a clear outlier. Next, in case of the averaged flux emitted in the backward direction (right hand side of 67),

<table>
<thead>
<tr>
<th>Intensity $10^{19}$ W/cm²</th>
<th>Temperature / MeV</th>
<th>Deviation δ/%</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>IP Lanex</td>
<td></td>
</tr>
<tr>
<td>3.66</td>
<td>0.38 0.41</td>
<td>-7.9</td>
</tr>
<tr>
<td>4.24</td>
<td>0.42 0.46</td>
<td>-9.5</td>
</tr>
<tr>
<td>4.39</td>
<td>0.62 0.54</td>
<td>12.9</td>
</tr>
<tr>
<td>4.55</td>
<td>0.47 0.52</td>
<td>-10.6</td>
</tr>
<tr>
<td>5.24</td>
<td>0.59 0.48</td>
<td>18.6</td>
</tr>
<tr>
<td>8.40</td>
<td>0.59 0.29</td>
<td>50.8</td>
</tr>
<tr>
<td>avg. rel. deviation δ:</td>
<td>6.52 0.55 0.76</td>
<td>-38.2</td>
</tr>
</tbody>
</table>

Table 2: Comparison of electron temperatures obtained by Image Plate and Lanex screen for both target normal directions. The temperature and the relative deviation between Lanex and IP temperature are given, and the average relative deviation for measurement series with $\geq 45$ shots. The data point for the single pulse event with $6.52 \times 10^{19} \text{ W/cm}^2$ is listed separately.
Figure 66: Example of an electron spectrum obtained with lanex (black plot) and IP (red plot).

one can see that there is a clear tendency towards larger yield when the laser intensity is increased. Again, the value at the $I \approx 6.5 \times 10^{19} \text{W/cm}^2$ is a single shot, only.

As regards the laser forward direction and the target parallel backward direction, results of the experimental intensity scans are shown in figure 68. In both cases, the total yield rises significantly as the intensity is increased. This general trend can be explained by an effectively larger spot size where the laser intensity is beyond the acceleration threshold. So a larger area of the target foil is irradiated with radiation-generating laser pulse intensity, which can produce more electrons.

In some geometries, practical limitations in the vacuum chamber, especially at reduced laser incidence angles, lead to the effect that the emission cones of target normal forward direction and laser direction overlap. In such cases, and when the electron beam in laser direction overlap, an image plate which is positioned in the target normal forward direction, may "see" the an electron bunch flux profile which s not centered on the image plate stack. Such a situation is depicted in figure 69. The figure further illustrates how flux data recorded on different stack IP’s is reconstructed. The first IP is located 5 cm behind the target foil (the radiation source), and is blocked by 1 mm Al to shield proton radiation. The optocouplers are located on the AL foil behind the from 1 mm Al shield (not shown in sketch). In the bottom part of the figure the data on the consecutive image plates, each separated by 1 mm of Al, is displayed. Naturally, the electron flux (red blue color scale) decreases from IP to IP due to the stopping of electrons. It is further seen that the hotter electrons are produced in the laser forward direction in this
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Figure 67: Total yield of electrons in target normal forward (left hand side of figure) and backward (right hand side) direction.

Figure 68: Total yield of electrons in laser forward direction and in target parallel backward direction.
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case, and drift our of the acceptance area of the IP stack when IP’s farther in the back of the stacks are regarded.

![Figure 69: Signal of electron flux on IP stack layers, where optocoupler shadows are present. In some geometries, the generated electron bunch emission cone as a mixture of target normal forward and laser forward direction processes cannot be caught completely by the IP stack.](image)

Large area IP stacks offer the possibility to examine the electron flux divergence in detail and energy dependent. Measured divergences on each stack IP for various scenarios are shown in figure 70. On the left hand side of figure 70 for the target normal forward direction, the divergence of each IP (where a lower number denotes an IP earlier in the stack, which therefore sees also lower energy electrons) is shown for four different cases: for an incidence angle of 45° at an intensity of $I \approx 4.4 \times 10^{19} \text{W/cm}^2$ (black plot) and at $I \approx 5.2 \times 10^{19} \text{W/cm}^2$ (red plot), for an incidence angle of 57.5° at an intensity of $I \approx 4.5 \times 10^{19} \text{W/cm}^2$ (green plot), and for an incidence angle of 70° at an intensity of $I \approx 3.6 \times 10^{19} \text{W/cm}^2$ (blue plot). A clear result is that the divergence of the higher energy electrons is lower when compared to the lower energy electrons. This is in accordance to references [8, 13, 104], where also a stronger collimation of higher energy electrons is seen. It is also very much in accordance with the underlying acceleration mechanisms and the figure-of-8 motion of electrons in the incident laser pulse field. As regards electron emission in the target normal backward direction, no such trend (if any, then in the opposite direction) was observed, which is in accordance with the dominant acceleration mechanisms and plasma expansion in this direction.

Using image plate stacks, the laterally resolved temperature distribution can be re-
Figure 70: Measured divergences of electron beams in the target normal forward and backward directions for 3 different angles at different intensities, obtained from IP stack measurements.

Figure 71 shows results of a number of spatially resolved measurements which confirm this characteristics of the electron bunches generated in laser-solid-interaction in our campaigns. This is in agreement with the divergence being larger for lower energies, compare the results displayed in 70. Shown in figure 71 a) are results for electron beams emitted in the target normal forward direction for three different angles of $45^\circ$, $57.5^\circ$ and $70^\circ$ in the beam center as well as in the beam rim (at the edge of the image plate). For these three laser incidence angles, figure 71 b) reveals that the temperature in the laser center is always higher than at the rim. It shall be noted that the presented data are single shots, because these experiments required image plate stacks without a central hole and without DUT’s and therefore not very well compatible with actual DUT irradiation shots.

Next, the electron spectra were optimized as regards their temperature, and averaged intensity scans have been carried through in order to be able to compare with predictions according to the scalings of Wilks and Beg, respectively. Figure 72 shows results of spectra measured with the magnetic spectrometers in target normal forward (a,b) and target normal backward (c,d) directions at different intensities. The resulting temperatures are in the range of $T_{\text{eff}} \approx 0.38$ to $T_{\text{eff}} \approx 0.95$ MeV in order to achieve the highest relevance for radiation belt electron spectral reproduction, see below. It is remarkable and a signature of the present acceleration scenario (prepulse, laser focus quality etc.) that the temperatures of the electrons emitted in the target normal backward direction (figure 72 a,b) were substantially larger when compared to those in the forward normal direction 72 c,d). These spectra were measured at the same time in both directions,
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which shows that scenarios can be found in which DUT’s can be tested with substantially different radiation spectral flux at the same time, which can be helpful if spectral flux scans are recommended for hardness assurance, for example for more complex missions on highly elliptical orbits where various flux levels and environments are passed due to the dependency of the spectral flux on distance to Earth, sun activity (AE8 model) etc.

![Graph](image)

Figure 71: temperature scalings.

Now these results are put into relation with the intensity scalings by Wilks [101]

\[ I = \frac{1.37 \times 10^{18} \text{W/cm}^2}{\lambda^2} \left( T_{\text{eff}}[\text{MeV}] / m_0 c^2 \right)^2 \left( 2 T_{\text{eff}}[\text{MeV}] / m_0 c^2 \right) = I[\text{W/cm}^2] \]

and Beg [102],

<table>
<thead>
<tr>
<th>Intensity $10^{19}$ W/cm$^2$</th>
<th>Spec. forward</th>
<th>Spec. backward</th>
<th>Temperature / MeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>1.2</td>
<td>--</td>
<td>--</td>
<td>1.02</td>
</tr>
<tr>
<td>1.5</td>
<td>--</td>
<td>--</td>
<td>0.68</td>
</tr>
<tr>
<td>1.6</td>
<td>--</td>
<td>--</td>
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<td>0.47</td>
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<td>0.96</td>
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<td>--</td>
</tr>
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<tr>
<td>8.4</td>
<td>0.59</td>
<td>0.88</td>
<td>1.02</td>
</tr>
</tbody>
</table>

Table 3: Overview of temperatures determined by magnetic spectrometers and image plate stacks for both target normal directions. The data in the forward normal emission direction are given separately for the central emission cone and the rim on the IP (where temperatures are lower).
Figure 72: Measured electron spectra for various intensities in target normal forward direction a) and b) and in target normal backward direction c) and d).
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Figure 73: Comparison of retrieved measured electron temperatures with the scalings of Wilks and Beg. a) shows the measured and retrieved temperatures of electron beams in the target normal forward direction obtained with permanent magnetic based spectrometers as well with IP stacks, and in b) the target normal backward direction temperatures obtained with a spectrometer and IP stack are depicted.
namely $T_{\text{eff}}[\text{MeV}]^3/\lambda^2 = I \times 10^{20} \text{W} \mu\text{m}^2/\text{cm}^2$. This comparison is shown in figure 73, where Wilks’ scaling (black line) and Beg’s scaling (red line) are plotted in comparison to measured data. In target normal forward direction (figure 73 a), stack data (green rectangle for beam rim, blue for center of beam) as well as permanent magnet spectrometer data (black rectangle) have been taken into account. In target normal backward direction, permanent magnet spectrometer data (black rectangle) as well as averaged stack data (green rectangle) have been regarded. In both cases, it is seen that the scaling of Beg fits much better than the Wilks scaling, which sheds light on the prepulse/preplasma situation during the actual runs and the dominant absorption mechanisms.

Figure 74: First laboratory based reproduction of space radiation (as present at GPS satellite level) with laser-plasma-generated bunches in the present project.)

Using the scaling and the comprehensive diagnostics, the electron temperatures could be further optimized and tuned to match the spectra flux present actually present in space. For example, figure 74 shows the spectral flux expected on GPS orbit at $L \approx 3.17$ (black solid line), an exponential fit (black dashed line) and optimized measured spectra in the target normal forward (red) and backward (green) direction. The left $y$-axis gives the electron flux per square centimeter, the right $y$-axis the electrons per MeV per msr. It is seen that there is an excellent agreement between the real flux in space and the lab radiation, especially in the important medium energy range. The lower energy range $< 1 \text{ MeV}$ is not exactly known since it is not very well measurable (this holds for both the situation in space as well as in the lab due to the spectrometer cutoff), neither is it particularly relevant (notable exception: surface charging) because low energies are more easily blocked. The mismatch in the higher electron energy range is also of secondary concern, because (note the logarithmic scaling) the number of electrons is much less here
in any case. Figure 75 shows a different scenario obtained during a different run to show that the space spectrum can be reproduced deliberately by tuning the parameters.

![Graph showing electron flux match of GPS level with a LPA-generated electron flux with $T_{\text{eff}} \approx 0.65$.]

This is one of the key results and main achievements of the present project. For the first time, "killer electron" space radiation spectral flux as present in the radiation belts has been produced in the laboratory here on Earth and has been used for irradiation of space electronics. This shows for the first time that it is possible to reproduce space radiation accurately in the laboratory, such that in the future it may no longer be necessary to work solely with monoenergetic particle beams.

The obtained particle spectra are furthermore very much in line with predictions coming from particle-in-cell simulations. Codes such as VORPAL as well as EPOCH have been used for this. Figure 76 shows simulation results obtained with EPOCH in 2D by Dr. Anupam Karmakar, an expert on laser-plasma interaction modelling, at the Simulation Laboratory Astro at the Leibniz Computing Center in Munich. It shows a laser pulse at $I \approx 4.5 \times 10^{19}$ W/cm$^2$ at Ti:Sapphire wavelength of $\lambda_L = 0.8 \mu$m and a duration of $\tau_L \approx 23$ fs focused to a spot size $w_0 \approx 3 \mu$m at $45^\circ$ incidence angle on a 30 micron thick Al foil and a tiny exponential ramp of 0.5 microns thickness due to preplasma formation. In the simulation the maximum electron density rises up to $30 \times n_e$.

On the left hand side of figure 76 the laser pulse is incident at $45^\circ$ incidence angle. The laser electric field (polarized in the simulation plane) is shown via the hot colorbar, and an interference pattern is visible between incident and reflected laser pulse at the beginning of the laser-plasma-interaction. Electrons are already seen (green) which penetrate the
target and will later leave the target on the back side after having propagated the 30 µm Al layer. The Al bulk electron density is shown with the reddish colorbar. The right hand side of the figure shows the spectrum of the electrons leaving the target in the normal direction after the interaction. The electron spectrum is clearly exponential, and a temperature fit results in $T_{\text{eff}} \approx 0.6$ MeV. This is in excellent agreement with the experimentally observed results, compare with figure 73 a).

Finally, let us have a look at the target normal surface parallel emission of electrons. Following the thesis of Königstein in the context of the present project [266] and [267], figure 77 illustrates the understanding of one of the crucial acceleration mechanisms which are present. This understanding was achieved in the context of the present project, showing vividly that the complex nature of laser-plasma-interactions still is a field of fundamental research. It shows that the laser pulse, incident under angle $\phi$, will generate a transient interference field which results as the interaction of incident and reflected laser pulse. During reflection, the laser pulse receives a phase jump, such that constructive interference fields are generated (indicated with the ellipses) which accelerate electrons rapidly into the target parallel direction(s).

The incidence angle defines the immediate acceleration direction, and other parameters such as plasma and laser pulse duration, power etc. substantially influence the final emission angle. An overview on the measured emission directions (making use of image plates and Lanex) in target "parallel" direction and relative intensities are shown in figure 78 for laser incidence angles of 45°, 52° and 65°. The emission angle is given both relative to the optical axis a) and to the bulk surface plane b). For these measurements, BK7 quartz glass wafers with a thickness of 4 mm and a surface roughness better than $\lambda/4$ were used as targets. Two main results are clearly seen in the figure. First, the
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Figure 77: One of the mechanisms at work for generation of acceleration of electrons in target parallel direction. A transient constructive interference resulting from incident and reflected laser pulse results in horizontal acceleration.

emission direction is closer to the target surface parallel direction when the laser pulse incidence is smaller, and second, the electron yield is larger when the incidence angle is higher.

Figure 78: Pointing and relative intensity of TSPA electrons in polar coordinates as observed during laser intensity and incidence angle scans.

In the runs with the thin Al foil and IP stacks and irradiation runs, the target surface parallel emission direction has been investigated in even more detail. As shown in figure 79, the available data indicates that a) there seems to be an intensity threshold in the target parallel backward direction after which the temperature of the emitted high energy electrons rises strongly. This intensity threshold in the campaigns was observed at approximately \( I \approx 5 \times 10^{19} \text{ W/cm}^2 \). In the target parallel forward direction, the temperature of the electrons is much higher on both the front (black) and back (red) side of the target, by a factor of more than two, see figure 79 b. The available data suggests that the temperature here rises linearly as intensity increases. Finally, figure 79 compares the temperatures observed on the target front side for incidence angles of
57.5° and 70°.

![Figure 79: Laser intensity scan and the resulting temperature of electrons emitted in the target parallel emission directions.](image)

The observed temperatures strongly suggest that the target parallel emission directions are highly relevant to space radiation testing. As regards the flux, figure 79 shows the electron yield per shot per square centimeter observed in the target forward parallel emission direction during intensity scans. Again, data suggests that there seems to be a threshold at intensity levels of a few $10^{19}$ W/cm², after which a jump in flux to very high levels is observed, and after which the yield continues to rise linearly. This is in agreement with our current understanding of the process, and demands further investigation – not only as regards space radiation reproduction, but on a very fundamental level for the whole laser-plasma interaction community.

The comprehensive diagnostics and thorough investigation and characterization of emission directions done within the present project allows for the first time a comprehensive view on the electron emission occurring during laser plasma interaction in the
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Figure 80: Laser intensity scan and the resulting yield of electrons emitted in the target forward parallel direction.

same campaign. Figures 81 and 82 show polar plots of the emission yield for the three angles 45°, 57.5° and 70° we focused on in our campaigns. Different intensities in the range of 3.7 to 6.5 × 10^{19} W/cm² are used. The laser pulse (red arrow) and the target angle (black line) are indicated, and the polar coordinate system follows our conventions as defined in figure 51. The directions start at 0° counterclockwise, such that ■ denotes the target parallel laser backward direction, ● denotes the target normal forward directions, ▲ the laser forward direction, ▼ denotes the target parallel laser forward direction and ♦ the target normal backward direction. One main message from this is that the target normal forward direction ● (the classical TNSA direction) as well as the novel TSPA direction ▼ are quite dominant, and therefore especially valuable for irradiation tests. The peak flux per shot obtained actually is reached for the TSPA direction ▼ at an intensity of 5.2 × 10^{19} W/cm², where ≈ 4 × 10⁹ electrons per cm² are measured in a distance of 5 cm to target. As always, variation of distance to target can increase or decrease this flux level, at the same time decreasing or increasing the irradiated area.

Figure 82 shows an analogous overview graph for the obtained temperatures, which are in the range from a few 0.1 MeV to a few MeV. The largest temperatures are reached in the laser forward direction ▲ and to a lesser extent in the target parallel forward TSPA direction ▼.

Based on the thoroughly examined and optimized electron flux reproduction and in close coordination with ESA, irradiation campaigns were carried through with the above mentioned optocouplers. Figure 83 shows on the left hand side the a few different optocouplers in front of a stack/spectrometer combination which were irradiated in the
Figure 81: Overview on spatial electron yield distribution per cm$^2$ (in 5 cm distance from target) for various angles of incidence and intensities.

(a) $I = 4.2 \times 10^{19}$ W/cm$^2$, 45°, 175 shots

(b) $I = 4.4 \times 10^{19}$ W/cm$^2$, 45°, 80 shots

(c) $I = 5.2 \times 10^{19}$ W/cm$^2$, 45°, 153 shots

(d) $I = 6.5 \times 10^{19}$ W/cm$^2$, 45°, 1 shot

(e) $I = 4.5 \times 10^{19}$ W/cm$^2$, 57.5°, 45 shots

(f) $I = 3.7 \times 10^{19}$ W/cm$^2$, 70°, 73 shots
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Figure 82: Overview on spatial temperature distribution for various angles of incidence and intensities.

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target normal forward direction. The right hand side of the figure shows the resulting radiography image of the corresponding image plate. This is a highly interesting picture in itself, as it reveals the inner structure of the optocouplers: one can see, for example, the pins and other inner details in the optocouplers. This demonstrates vividly the extremely high accuracy of our system. A high level of detail can also be seen on the other stack image plates – this is attributed to the extremely high laminarity of electrons (and protons) from laser-solid interaction. The green dashed rectangle next to one of the optocouplers on the IP (red dashed rectangle) is used to determine the total irradiation level. In this case, in this rectangle \( \approx 2 \times 10^9 \) electrons were detected by the IP. This low electron number is already enough to resolve the depicted detailed structure – which is another advantage of using the accurately calibrated and highly dynamic IP electron response as flux monitors.

![Figure 83: Left: Set of optocouplers on mounting foil. Right: Radiography image on image plate resulting from irradiation during a single shot.](image)

To determine the performance of the optocoupler DUT’s before and after irradiation, the Current Transfer Ratio between output and input current \( \text{CTR} = \frac{I_c}{I_f} \) is determined. For this, we have used specialized equipment at ESTEC, which is shown in figure 84. At ESTEC, the Agilent Test Fixture 16442A, which provides various sockets (DIL, TO-5), was used in combination with the precision semiconductor parameter analyzer Agilent 4156C, which can provide currents down to 1 pA. The climatized test environment at ESTEC allows to rule out temperature effects on optocouplers during CTR determination, which can have significant effects on the semiconductor behaviour.

A clear irradiation damage is shown in figure 85 for the optocoupler model Vishay
SFH6345, for which the CTR was measured for input currents of $I_f = 100 \mu A$ and $I_f = 1 \text{ mA}$, respectively, before and after irradiation. In figure 85, the CTR of not irradiated (but otherwise treated and stored in the same conditions as the other optocouplers) reference optocouplers (DUT’s # 30-60) is displayed and encircled with a blue dashed ellipse. As is clearly seen, the CTR of the not irradiated optocouplers is unchanged. In this campaign, there was one group of optocouplers (tagged with ”1” and ”2” in figure 85, respectively) which was exposed to $3.2 \times 10^6 \text{ e}^-/\text{cm}^2$ in the target normal forward and backward direction, respectively, in one irradiation block and then in combination with a second irradiation block with in total $4.6 \times 10^7 \text{ e}^-/\text{cm}^2$. After exposure to $3.2 \times 10^6 \text{ e}^-/\text{cm}^2$ the optocouplers were again taken from Düsseldorf to Noordwijk to ESTEC, and the CTR was determined again. No significant degradation of CTR was observed. Even after the next irradiation block in Düsseldorf, when the fluence was increased to $4.6 \times 10^7 \text{ e}^-/\text{cm}^2$, hardly a degradation of CTR performance could be measured at ESTEC. In contrast, optocouplers from group 3 and 4, which were exposed to a far higher fluence, namely at maximum $2.1 \times 10^9 \text{ e}^-/\text{cm}^2$, show a significant (a thorough error analysis was performed) and cumulative deterioration after each irradiation block. The maximum CTR degradation was $> 3\%$. This example of successful deterioration of performance is one core result of the present campaign. It was shown that using accurately reproduced space radiation in the laboratory with laser-plasma accelerators, testing is possible making use of adapted standard testing techniques and that significant radiation damage can be exerted on DUTs with these laser-produced radiation belt electrons.

In the proof-of-concept experiments, the focus was put on accurate characterization, monitoring and optimization of the electron flux. The focus was not on application of maximized fluence, which is straightforward. In the campaigns carried through in the context of the present project, a number of factors limit the averaged flux and the total...
Figure 85: DUT degradation after irradiation: Optocoupler CTR degradation after irradiation with laser-plasma-produced radiation belt flux.
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fluence:

1. the time needed to break the vacuum to retrieve image plates, to insert new image plates and to evacuate the chamber (approx. 1 hour per irradiation block)

2. the time needed for the image plate readout and erase processes (approx. 20 minutes)

3. the time needed to bring the target foil into the Raleigh length of the strongly focused laser pulse (hour-scale)

4. the limited number of shots applicable on one target foil, and the time needed to introduce a new target foil (tens to hundred of shots per foil)

5. the pump performance needed to keep the chamber evacuated during shots

6. the repetition rate of the laser system (currently 10 Hz)

All of these limitations can be overcome already with today’s state-of-the-art technique. This means that the number of shots per time interval, and therefore the averaged flux, can be relatively straightforwardly increased by many orders of magnitude. First, the vacuum does not need to be broken for image plate change because one can rely on the cross-salibrated lanex response for online monitoring, which is thorough enough. Image plates inside the chamber can still be used but need to be more heavily shielded and can then nevertheless provide useful fluence information once the irradiation has ended. Then, the image plate readout process time is consequently also not relevant anymore. Next, positioning of the target foil is much easier when a softer focusing is used. While currently, the Raleigh length $z_R = \frac{\pi w_0^2}{\lambda} = 30 \text{ µm}$ due to the very strong focusing with the F/2 parabola, a longer Raleigh length to values beyond 100 microns will dramatically relax the demands put on target positioning. It will furthermore stabilize the radiation output because variations in position do have much less effect on the laser-plasma-interaction. In this connection, and with regard to the maximum number of shots on one target foil, this can be overcome with tape drives. Such tape drives, e.g. consisting of tens of meter long VHS video band, for example, are well-known tools in the laser-plasma-community [268–270]. We have begun construction of a tape drive which is suitable for space radiation reproduction in the context of the present project (M. Quast et al. at University of Hamburg). An alternative to this if laser-solid-interaction is chosen as underlying acceleration mechanism would be droplet targets, such as described in [271], for example. If laser-underdense interaction is chosen as acceleration mechanism, then a steady-state gas cell with differential pumping would solve the problem of repetition rate. Also, reduced debris as with droplets [271] and underdense targets would substantially decrease the requirements put on the vacuum system and the pumps. Finally, as regards repetition rate, already today kHz system
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with many mJ of laser energy and pulse durations < 100 fs are commercially available. The rapid advance in laser technology has already even produced 100 kHz-level laser systems based on OPA, fiber lasers and thin disk lasers. This trend, and furthermore a much better wall plug efficiency with fiber lasers etc. will continue for the foreseeable future.

It is therefore reasonable to estimate the irradiation time needed with laser-plasma-systems for elevated repetition rates of 10 Hz and belong, which would be seen by future projects, and based on the flux per shot we have seen in the proof-of-concept experiments. Figure 86 shows the irradiation times which would be needed to reproduce the flux levels encountered on a typical navigation orbit, which amounts to $\approx 3 \times 10^{12} \text{e}^-/\text{cm}^2$. This shows that various emission direction can be used to deposit the fluence encountered on electronics onboard a satellite in under 1000 seconds.

![Figure 86: Irradiation time required to produce the daily fluence on Nav orbit of $\approx 3 \times 10^{12} \text{e}^-/\text{cm}^2$ in different emission directions at 1 kHz repetition rate.](image)

To put this into perspective and as further outlook, figure 87 visualizes a comparison between the fluence level possible obtainable from laser-plasma-accelerators in various directions at 10-100 kHz when compared to linacs, one assuming a flux of $\approx 1.3 \times 10^8 \text{e}^-/\text{cm}^2$ (the low flux linac, orange bar in the figure), and one assuming a flux of $\approx 1.2 \times 10^{10} \text{e}^-/\text{cm}^2$ (the high flux linac, orange bar in the figure). The low flux linac

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level of fluence is approximately reached already with a 10 Hz laser system, while at 10 or 100 kHz the fluence from laser-plasma-accelerators would surpass the fluence level even of high flux linacs by orders of magnitude. This is an important feature, since generally one wants to keep irradiation times as small as possible. It shall be noted that at the same time the spectral flux of the laser-plasma-accelerator would be broad-band and space like. This clearly demonstrates the potential of laser-plasma-accelerators for space radiation testing on various levels.

Figure 87: Fluence levels reached with laser-plasma-accelerators when compared to low and high flux linacs after 1000 s.

Finally, another bar diagram is depicted in figure 88 which is complementary to figure 87 in showing the time DUT’s would need in the different LPA emission directions when compared to low and high flux linacs to reach a certain fluence level, again the value of \( \approx 3 \times 10^{12} \text{e}^-/\text{cm}^2 \) which is reached at an orbit of \( L = 3.5 \) during maximum sun activity (AE\textsubscript{max}) important for GPS satellites and the like. Here, it is seen that the time required to reach these flux levels can be substantially lower than with the linacs.
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Figure 88: Irradiation time required to produce the daily fluence on Nav orbit of \( \approx 3 \times 10^{12} \text{e}^-/\text{cm}^2 \) in different emission directions at different rep rates in comparison with low and high flux linacs.

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10 Summary and Outlook

Laser-plasma accelerators (LPA’s) have been used for the first time for accurate reproduction of space radiation and for successful controlled testing of performance degradation of electronic components. LPA is a relatively young, but fast rising technology which brings various kinds of accelerator R&D into the single university lab scale. This is done by the giant fields which are available in laser-plasma-interaction, which are orders of magnitude higher then in conventional state-of-the-art accelerators. In turn, the accelerators are shrunk down in size, because no longer accelerator distances of tens or hundreds of meter are necessary anymore. Second, both understanding and technology readiness level of laser-plasma-concepts as well as laser technology itself has made much progress in past years and continues to do so. The time seems right to establish laser-plasma-acceleration for space radiation testing. The laser-plasma-community has demonstrated the production of quasi-monoenergetic electrons with energies up to more than 4 GeV, and proton beams up to a few hundreds of MeV. As a matter of principle, producing high energy electrons is much easier than producing high energy protons. At the energy frontier, one needs generally the biggest lasers and will only achieve limited stability in particle bunch generation. While this frontier is constantly pushed and the stabilization is increased, it is much easier to generate low energy electrons and protons, in the range of a few MeV. The laser systems needed for this are much more cost-effective, smaller and can have higher repetition rates, which increases the obtainable flux. Moreover, producing broadband spectral flux is much easier with laser-plasma-accelerators when compared to monoenergetic beams. This stands in diametral contrast to conventional linacs and/or cyclotrons, where the beam output is monoenergetic. It is well know that practically all kinds of space radiation have broadband spectral flux. This is the inherent regime of laser-plasma-accelerators. It is much easier to produce broadband flux of a few MeV electrons than to produce high energy, monoenergetic electrons. So called ”killer electrons” in the Earth’s radiation belt are especially prominent at GPS levels. They have broadband spectra and energies up to a few MeV – perfect candidates for proof-of-concept measurements.

Laser-plasma-acceleration based on laser-solid interaction was carried through in various campaigns at the Arcturus laser system at University of Dusseldorf. A milestone was reached in that for the first time, killer electron flux from space has been very accurately reproduced. This was achievable due to the large tunability of laser-plasma accelerators and well established engineering scaling laws which give the user various knobs to turn to control the radiation output. By turning these knobs, and by employing sophisticated diagnostics, the spectral flux output could be designed to match those on various orbits and environments in space, which were determined using the AE8 model. The diagnostics used are actually hybrids, having partially been adapted from the conventional accelerator community, but to an increasing part diagnostics developed in the laser-plasma-community are finding their entrance into the conventional accel-
This shows how research in both fields increasingly cross-fertilizes each other, and at various fronts the both communities are merging. Having optimized the output from laser-plasma accelerators, we then have adapted a test procedure from the radiation testing community and have introduced optocoupler test devices into the laser-plasma-generated space radiation in the target chamber. By applying this "space radiation reproduced in the lab on Earth" at different fluence levels, we could demonstrate that this type of space radiation is useful for the space radiation field.

Hardly such a breakthrough could come at a more propitious time, since on the one hand there is an increasing shortness of beamtime for radiation testing due to the increasing demand and also due to the increasing complexity of space electronics, and on the other hand space-radiation capable laser-plasma accelerators are mushrooming all over the member states. There is the opportunity to establish laser-plasma accelerators as complementary radiation sources which could then support the established radiation sources. Laser-plasma accelerators could be used for specialized tasks such as accurate reproduction of certain kinds of space radiation, lower energy and broadband electrons and lower energy protons and ions, for example. It is intriguing that laser-plasma accelerators can produce electrons as well as protons and ions in a wide range of parameters, and also it is possibly to switch from electron generation to proton generation with minimal change to the setup – in fact, it is even possible to produce electrons and protons at the same time.

Two basic types of plasma acceleration are distinguished, one where the laser pulse interacts with overdense material (such as target foils), and the other where the laser pulse interacts with underdense, gaseous material. While both types of acceleration are good to produce high flux of electrons, here we focused on laser-solid interaction and produced and diagnosed various emission directions around the target. A focus was put on diagnosis and characterization of this radiation. All these emission directions are useful for space radiation testing. The detailed diagnosis required to break the vacuum to take out imaging plates which were used as highly accurate diagnostics, which resulted in interrupts of the irradiation. These interrupts and others due to the target holder etc. limited the obtainable average flux. However, at 10 Hz repetition rate or at even higher repetition rates such as 1 kHz (state-of-the-art) and 100 kHz (prototypes exist), extrapolation shows that the flux is extremely high and can exceed that of conventional linacs by a wide margin. The daily electron flux impinging on satellite electronics situated in the radiation belt may be reproducible in less than a minutes’ time.

Future work should include the increase of the repetition rate, e.g. by using tape-drives or underdense targets, as well as on harvesting the low energy protons, as well as effects on high peak flux, which is obtainable close to the target and is otherwise not accessible. Various additional test procedures have to be implemented and diagnostics have to be be adapted. Generally, the radiation effect community and the laser-plasma community should be developing this technique in a joint programme. Such a programme would contribute to position the space community in the member states in a spearhead-
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ing position for the future exploitation of laser-plasma-accelerators. Another topic of future research should be hybrid systems, where a plasma section is used to broaden the spectral flux from linacs [272]. Dedicated beamlines at large laser-plasma-facilities such as the Scottish Centre for the Application of Plasma-based Accelerators SCAPA or sites belonging to the European Extreme Lights Infrastructure may be an efficient and continuity-safe way to go forward.
11 Acknowledgements

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