Doctoral Thesis

Energetic particles and plasma instabilities in solar flares

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Energetic Particles
and Plasma Instabilities in
Solar Flares

A dissertation submitted to the
SWISS FEDERAL INSTITUTE OF TECHNOLOGY
ZÜRICH

for the degree of
Doctor of Natural Sciences

presented by

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2002
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Abstract

Many astrophysical sources including our Sun exhibit signatures of non-thermal particles which manifest themselves in emissions at various wavelengths. While those are the only source of information in most astrophysical plasmas, the Sun allows us to incorporate \textit{in situ} measurements of the solar wind to the analysis. Moreover, due to its proximity many features can be spatially resolved. Solar flares therefore offer the most comprehensive source of information on particle acceleration and provide a link between astrophysical and laboratory plasmas. These observations in conjunction with theoretical modeling may reveal the basic mechanisms of particle acceleration on the Sun, and by extension also in other astrophysical plasmas.

The work presented here starts with the analysis of radio observations of the Sun. So-called type III bursts, radio signatures from escaping electron beams on ‘open’ field lines, associated with narrow band metric spike emission have been analyzed to corroborate the interpretation of dynamic spectra in terms of signatures of the particle acceleration region. The radio spike emission is thought to be closely related to the acceleration mechanism and should therefore spatially coincide with the origin of the trajectory of the type III producing electrons. It was found that in all of the analyzed events the spike sources are always located at positions coinciding with the expected locations obtained from extrapolating the electron trajectories to lower altitudes. In one of the events, two subsequent type III bursts originated in the same source region allowing us to determine the acceleration region very accurately by intersecting the two trajectories. The findings strongly support the generic flare picture, where the energy release takes place in the corona above magnetic loop structures.

While this analysis yields insight into the spatial location of the acceleration region, the nature and properties of the accelerating process are addressed in the rest of the work. Candidates for the accelerating process can roughly be divided into three broad classes: \textit{parallel DC electric fields, shocks} or \textit{stochastic acceleration by plasma waves}. The different aspects of each class have not been addressed in detail in this work. A broad approach was chosen instead by trying to extract the main feature(s) common to all members of the above classes. This allows us to derive general plasma properties without choosing a specific scenario. The main common feature that was found is an anisotropy $T_\perp < T_\parallel$ that builds up in the course of the acceleration process in the velocity distribution function of the electrons ($\parallel$ and $\perp$ here refer to the directions relative to the background magnetic field). Such an alignment along the background magnetic field may cause the acceleration to cease. As a consequence, many of the aforementioned accelerators need a pitch-angle scattering process that erodes the anisotropy by

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redistributing the particle energies isotropically.

A linearized kinetic approach was chosen for a stability analysis of such an anisotropic electron population to plasma waves. It was established that the Electron Firehose instability (EFI) has to be expected during solar flares. Left-hand circularly polarized electromagnetic waves with frequencies around the proton gyrofrequency (EF waves) are non-resonantly excited by the electrons. The response of two different plasma components to the excited waves was investigated: i) The kinetic response of the electron population to the instability and ii) the acceleration and enrichment of $^3$He have been analyzed analytically and by test particle simulations. It was found that the electrons experience a translation to higher pitch-angles in velocity space in the presence of a field of EF waves. The EFI is therefore a viable agent for the required pitch-angle scattering. From the latter work it is learned that $^3$He can be readily accelerated by EF waves to the required energies or even beyond. Moreover, the spectral properties of the EF wave spectrum result in an enhancement of the abundance of $^3$He above $^4$He, reproducing the observations of various spacecraft.

The EFI as an inherent component of bulk electron acceleration provides therefore two major ingredients to flare particle acceleration: it can pitch-angle scatter electrons to higher perpendicular velocities and therefore enable efficient acceleration and it is able to accelerate and enrich $^3$He via EF waves. Both aspects are explained as an inherent feature of the acceleration process of electrons in solar flares.
Zusammenfassung


Während dieser Teil der Arbeit die räumliche Position der Beschleunigungsregion behandelt werden im Rest der Abhandlung die die Eigenschaften des Beschleunigungsprozesses untersucht. Mögliche Kandidaten für den Beschleuniger können grob in drei breite Kategorien unterteilt werden: Parallele elektrische Felder, Schocks und stochastische Beschleunigung durch Plasma Wellen. In dieser Arbeit wurden nicht die spezi-
Zusammenfassung


Um die Stabilität einer solchen Elektronenverteilungsfunktion gegenüber Plasmawellen zu untersuchen, wurde ein Zugang in linearisierter kinetischer Theorie gewählt. Es konnte gezeigt werden, dass das Auftreten der Elektronen Firehose Instabilität (EFI) während solaren Flares zu erwarten ist. Linkshändig zirkular polarisierte elektromagnetische Wellen mit Frequenzen nahe der Protonen Zyklotronfrequenz (EF Wellen) werden auf nicht-resonante Weise von den Elektronen angeregt. Weiterhin wurde die Reaktion zweier Plasmakomponenten auf die erzeugten Wellen untersucht: $i)$ Die kinetische Antwort der Elektronen auf die Instabilität und $ii)$ die Beschleunigung und Anreicherung von $^3$He wurden vermittels analytischer Rechnungen und Testteilchensimulationen untersucht. Im ersten Teil der Untersuchung wurde gefunden, dass die Gegenwart eines EF Wellenfeldes die Elektronen im Geschwindigkeitsraum zu höheren Pitch-angles versetzt. Die EFI ist folglich ein guter Kandidat für die erforderliche Pitch-angle Streuung. Die zweite Untersuchung ergab, dass $^3$He durch EF Wellen direkt zu den beobachteten Energien und sogar weiter beschleunigt werden kann. Darüberhinaus ergeben die spektralen Eigenschaften des EF Spektrums eine erhöhte Anreicherung von $^3$He gegenüber $^4$He, was auch den Beobachtungen verschiedener Satelliten entspricht.

Die EFI als inhärenter Teil der Bulk Beschleunigung von Elektronen liefert folglich zwei wichtige Bestandteile für die Teilchenbeschleunigung in solaren Flares: Sie kann Elektronen zu grösseren Pitch-angles streuen, was effiziente Beschleunigung ermöglicht und vermag $^3$He vermittels EF Wellen zu beschleunigen und anzureichern. Beide Aspekte werden als intrinsischer Bestandteil des Beschleunigungsprozesses von Elektronen in solaren Flares erklärt.

$^4$Pitch-angle – Anstellwinkel des Teilchens im Geschwindigkeitsraum
Chapter 1

Introduction

felix qui potuit rerum cognoscere causas
Publius Vergilius Maro (70 BC–19 BC)

1.1 Solar Flares

Flares of the Sun are among the most energetic phenomena in the solar system. They come in a wide variety of spatial extent, temporal behavior, spectral shape and energy output. Energies released range somewhere between $10^{28} - 10^{34}$ ergs and vary on time scales from fractions of seconds to several tens of minutes (e.g. Švestka, 1976; Kiplinger et al., 1984). For comparison: this amount of energy $\sim 10^{32}$ ergs is about large enough to lift the moon to a $50$ km higher orbit. Much of the energy first resides in supra-thermal particles that either remain trapped at the Sun and produce a variety of radiative signatures (Ramaty & Murphy, 1987) or escape into interplanetary space (Reames, 1990) and propagate to Earth and beyond. Radiative signatures of flares comprise continuum emission spanning the range from radio and microwave wavelengths to soft X-rays (SXR) ($\sim 1 - 10$ keV), hard X-rays (HXR) ($\sim 10 - 300$ keV), and finally gamma-rays (above $\geq 300$ keV), as well as narrow gamma-ray nuclear de-excitation lines between $\sim 4 - 8$ MeV. Solar flares thus offer a wide diversity of diagnostics for the particle acceleration mechanism(s), to understand which is the ultimate goal of flare research.

On the basis of the duration of their SXR emission, solar flares can be subdivided into two main classes: gradual and impulsive events (Pallavicini et al., 1977). Impulsive events produce short duration SXR and HXR radiation. Also, they tend to be compact and to occur low in the corona, while gradual events occur at greater heights and exhibit long-duration X-ray emission. Besides the time duration and morphology, the classes are distinguished by their variation in the abundances of energetic particles (see reviews by Lin, 1987; Reames, 1990). Gradual events generally exhibit large proton fluxes and produce energetic ions with coronal abundances. They are often associated with Coronal Mass Ejections (CME) and/or type II radio emission. Impulsive events, in contrast, produce ion abundances that dramatically deviate from coronal values (Reames et al.,
Chapter 1. Introduction

Figure 1.1: Radio spectrogram of a large group of type III emissions with clusters of associated metric spikes at higher frequencies. This event is further analyzed in Chapter 2 and corresponds to the spectrum displayed in Fig. 2.2.

1994). Due to the extraordinarily high $^3\text{He}$ enhancement during such events (up to a factor of $\sim 10^4$), they are also referred to as $^3\text{He}$-rich events. Impulsive events are associated with $\sim 2 - 100$ keV electrons that produce type III radio signatures (see Fig. 1.1).

The standard model for these observations describes energetic particles in gradual events as the result of a CME-driven shock, while the energetic particles in impulsive flares are accelerated by another yet unknown mechanism or mechanisms. Refinements of the two-class picture (e.g. Mandzhavidze & Ramaty, 1993; Cliver, 1996) suggest, however, that it is the same mechanism in impulsive and gradual events that accelerates the particles which remain trapped at the Sun, and that escaping particles observed in interplanetary space are similar to these trapped particles. In other words, there is an impulsive core in both classes of events that produces the trapped particles which escape, as in pure impulsive events, and may join the CME-shock accelerated particles, as in gradual events. In the following, the word ‘flares’ is usually used to refer to the impulsive class of events although, due to the above, the discussion can also be true for the not-shock accelerated fraction of energetic particles in gradual events. A summary of the two class picture of solar flares is shown in Table 1.1.

The present canonical flare picture (see Fig. 1.2) describes ion and electron acceleration as a result of restructuring magnetic fields high in the corona. The overall magnetic structure contains loop-like field lines, i.e. field lines being anchored at both ends in the


<table>
<thead>
<tr>
<th></th>
<th>Impulsive</th>
<th>Gradual</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^3$He/$^4$He</td>
<td>$\sim 1$</td>
<td>$\sim 0.0005$</td>
</tr>
<tr>
<td>Fe/O</td>
<td>$\sim 1$</td>
<td>$\sim 0.1$</td>
</tr>
<tr>
<td>H/He</td>
<td>$\sim 10$</td>
<td>$\sim 100$</td>
</tr>
<tr>
<td>$Q_{Fe}$</td>
<td>$\sim 20$</td>
<td>$\sim 14$</td>
</tr>
<tr>
<td>Duration</td>
<td>up to Hours</td>
<td>up to Days</td>
</tr>
<tr>
<td>Longitude Cone</td>
<td>$&lt; 30^\circ$</td>
<td>$\sim 180^\circ$</td>
</tr>
<tr>
<td>Radio Type</td>
<td>III,V,(II)</td>
<td>II,IV</td>
</tr>
<tr>
<td>X-Rays</td>
<td>Impulsive</td>
<td>Gradual</td>
</tr>
<tr>
<td>Corona</td>
<td>–</td>
<td>Coronal Mass Ejection</td>
</tr>
<tr>
<td>Solar Wind</td>
<td>–</td>
<td>Interplanetary Shock</td>
</tr>
<tr>
<td>Events/Year</td>
<td>$\gtrsim 10000$</td>
<td>$\sim 10$</td>
</tr>
</tbody>
</table>

**Table 1.1:** Properties of impulsive and gradual solar flares (after Cliver, 1996).

lower chromosphere and photosphere, and additional 'open' field lines that are anchored only at one end while the other end extends into interplanetary space and returns to the Sun after an immensely longer path.

Magnetic reconnection, as the primary energy release, is thought to trigger the conversion of magnetic energy into bulk mass motion or out-flowing energetic particles. The reconnection of magnetic field lines is a result of driving oppositely directed fields together. The outflow velocities of the reconnection site can exceed the local Alfvén speed and therefore excite turbulence (region A in Fig. 1.2, see also Section 1.3.3). The expectation of finding the particle acceleration site above the loop are supported by observations of the Yohkoh spacecraft published by Masuda et al. (1994). HXR sources above the loop top have been observed. Such a source is difficult to explain by assuming the primary acceleration site to lie somewhere else than above the loop, e.g. at the loop foot points. In addition, observations of so-called radio spike emission at metric wavelengths also indicate that the accelerator is located above the loop (Krucker et al., 1995; Paesold et al., 2001; Benz et al., 2002, see also Chapter 2). This geometry naturally leads to trapped (closed field lines) and escaping (open field lines) energetic particles. Only above the loop are open and closed field lines directly connected, thus permitting a simultaneous injection of energetic particles.

It should be mentioned here, that the actual accelerator itself seems to be hidden from the observer. The only candidates for offering a first-hand view of the acceleration process could be the aforementioned radio spikes: narrow band ($\sim 5 \%$ of the central frequency) radio emissions at either metric or decimetric wavelengths on time scales of $\leq 100 \text{ ms}$. An example of metric spike emission is shown in Fig. 1.1 (see also Chapter 2).

After being accelerated by some mechanism(s) left unspecified here, energetic elec-
Figure 1.2: Sketch of the present canonical flare scenario. A) Initial energy release: Magnetic reconnection. B) Particle acceleration site: Radio spike emission and HXR emission on loop-top. C) Downward-traveling energetic electrons emit gyro synchrotron radiation and may emit reverse-drifting dm type III radio emission. Accelerated electrons impact the chromosphere, causing HXR emission via thick target bremsstrahlung. In addition, gamma-ray line emission comes from the foot points. Chromospheric material is heated, resulting in evaporation into the loop. The filled loop radiates in EUV and SXR. D) Accelerated electrons escaping on open field lines to interplanetary space can form beams and cause meter wave type III radio emission.

Electrons stream along magnetic field lines and emit, at relativistic speeds, gyro synchrotron radiation at microwave wavelengths. Also, the energetic electrons can build up beams that may excite Langmuir waves which eventually are converted via non-linear wave-wave scattering to observable radio waves, giving signatures of reverse-drifting decimetric type III emission.

Later, when arriving at the loop foot points, the electrons penetrate the chromosphere (Fig. 1.2C). The interaction with ambient protons produces HXR emission. However, far more and most of the kinetic electron energy (∼99.99%) goes into heating the ambient cooler electrons in the chromosphere via Coulomb collisions. Since nearly all of the super thermal energy in the electrons is lost to the ambient electrons, the resulting emission is referred to as thick-target non-thermal bremsstrahlung and can be observed in the HXR range.

The heating of chromospheric material results in the so-called ‘chromospheric evaporation’ (e.g. Mariska et al., 1993): The heated chromospheric material expands upwards into the corona following the guiding field lines and filling the loop structure with hot
material. These flows emit relatively long-lived thermal extreme ultraviolet (EUV) and SXR radiation (see also Fig. 1.3).

Energetic electrons on open field lines can radiate at radio wavelengths in the metric range by virtue of the same mechanism as downward-traveling electron beams, and so produce classic type III radio signatures (see Fig. 1.1). When propagating further out, the beams can be followed through interplanetary space as interplanetary type III radio emission (e.g. Aschwanden & Benz, 1997) and the beam electrons themselves can be observed by spacecraft (e.g. ISEE 1–3, WIND or Advanced Composition Explorer (ACE)) at 1 AU (e.g. Krucker et al., 1999; Lin et al., 1981). Observability of interplanetary solar energetic particles requires magnetic connectivity of the spacecraft with the source region and is therefore very selective.

Energetic ions either escape on open field lines and can be observed by in situ measurements in space (e.g. for $^3$He observations Shaeffer & Zahringer, 1962; Hsieh & Simpson, 1970) or interact with ambient nuclei in the relatively dense chromosphere (for reviews see e.g. Chupp, 1984; Ramaty & Murphy, 1987). This results in excited nuclei, radioactive nuclei, pions and neutrons. Narrow de-excitation lines are a result of interactions between H and $^4$He nuclei with energies of $\approx 1$ and $\approx 100$ MeV amu$^{-1}$ on the one hand and ambient heavier nuclei. The inverse process, i.e. energetic heavier ions impacting ambient H and $^4$He yields broader lines of about $\sim 1$ MeV widths. An

**Figure 1.3:** This image of coronal loops over the eastern limb of the Sun was taken in the TRACE 171 Å pass band, characteristic of plasma at $10^6$ K, on 99/11/06. The image is rotated by 90 degrees.
additional deuterium formation line is observed at $\sim 2.2$ MeV. More information on
the accelerated ion population can be obtained by modeling pion decay emission which
is a secondary effect of high energy protons (which dominate pion production).

The preceding discussion shows the wide diversity of diagnostics that are available
for solar flare research. However, it will not be possible to have in situ measurements of
solar flares that would allow us to probe a theoretical model by witnessing first-hand the
effective processes. Therefore, the only goal of flare modeling is a close description of
the greatest amount of solar flare data and its stability to initial conditions. Condensed
observational requirements for electrons and ions are formulated in the next section to
eliminate singular events and to find the main observational evidence that a successful
model of particle acceleration has to account for.

1.2 Central Observations of Energetic Particles

A summary of the essential observational requirements for a successful particle accel¬
eration mechanism is presented. Most of the information on the number of interacting
(trapped) electrons and on their acceleration time scales are obtained by analyzing
HXR emission. The required information can be extracted from their spectra by fit¬
ting bremsstrahlung spectra to the HXR profiles. In addition, line emission of nuclear
de-excitation lines and pion radiation can be modeled to gain information on the in¬
teracting (trapped) ions. Information on ion abundances is gained directly through in
situ observations in space. It is important to note that the features presented below are
not all found in every individual flare. They rather represent averaged aspects of flare
particle acceleration that have to be accounted for by a viable acceleration model.

1.2.1 Electron Acceleration

Bremsstrahlung photons of a given energy are produced principally by electrons of com¬
parable energies. It can therefore be directly concluded from HXR emissions during
solar flares that electron energies range up to relativistic energies $\sim 100$ keV (Dennis,
1988). Gamma ray emission from electron dominated events also indicates the presence
of ultra-relativistic electrons at energies of tens of MeV. The time scales for the acceler¬
ation up to $\sim 100$ keV can be estimated from the temporal variation of HXR emission.
Observations by e.g. Kiplinger et al. (1984) and Machado et al. (1993) revealed fine scale
structures on time scales of $\sim 400$ ms in the HXR emission of impulsive solar flares.
Since HXR emission is a signature of accelerated electrons, the observed time scales also
correspond to the electron acceleration time scales. In general, electrons are accelerated
to energies of $\sim 100$ keV on time scales of the order of $\sim 1$ s. Energies above these
values are reached on longer time scales. The analysis of time profiles of gamma ray
emission during electron dominated gradual events by Rieger (1994) yields typical time
scales for these emissions from a few second to about 30 s.

While the above describes the highest observed energies that are naturally only
reached by a few particles, the bulk of the electron population is energized to about
~ 10 keV. 'Energization' here refers to a quasi-heating process that maintains an isotropic state of the particle distribution function throughout the acceleration process. Time integrating the HXR profile of impulsive flares yields a total energy content in the bulk electron component of about $10^{31}$ ergs. Energization rates of about $\approx 10^{36} - 10^{37}$ electrons s$^{-1}$ are required to account for the observed energies. A typical flaring volume of $\sim 10^{27}$ cm$^3$ therefore has to be depleted within about $\sim 1$ s when assuming a coronal electron density of $\sim 10^{10}$ cm$^{-3}$. This poses a very stringent so-called 'number requirement' to the accelerator. In order to process such a large number of electrons on the required time scales, the accelerator must be very efficient and even real-time replenishment of the source region has to be provided. This is a critical item a possible accelerator must be consistent with.

Analysis of the time profiles and bandwidths of electromagnetic radiation from solar flares as in radio (Benz, 1985) and HXR (Kiplinger et al., 1984; Machado et al., 1993) emission further indicate a fragmentation of the energy release into multiple smaller events. Typical time scales derived from X-ray observations are around $\sim 400$ ms. The number requirement therefore can be reformulated: Electron acceleration is split into energy release fragments (ERF) requiring electrons to be accelerated up to energies of 20 keV at a rate of $\sim 5 \cdot 10^{34}$ electrons s$^{-1}$ for about $\sim 400$ ms. Each of the ERF's then contributes about $\sim 5 \cdot 10^{26}$ ergs to the total energy content in the bulk electrons. The ERF's are apparent only in moderate flares, where the actual number of them at any time is sufficiently small to be observed separately. In larger events they blend together to form a smoother HXR profile.

### 1.2.2 Ion Acceleration

Trapped ions on the Sun are briefly addressed here while more detail is given to ion abundances. Observed from space, the escaping ion populations of a flare and their abundances are a very important diagnostic for probing the accelerator.

The main requirements for the ion accelerator in an impulsive flare can be summarized in two major observational evidences: i) The time scale for acceleration of ions to energies of about $\sim 100$ MeV is of the order of $\sim 1$ s while the higher energies observed (up to $\sim 1$ GeV) are reached on time scales of a few seconds. ii) The rates at which protons, as the major ion component, have to be accelerated to more than approximately $\sim 1$ MeV is about $\sim 10^{35}$ s$^{-1}$ for the duration of several seconds. The total energy content is about $\sim 10^{31}$ ergs. Since these data are obtained from modeling gamma ray emission line profiles, the ion energies involved are very high and the observed population therefore supra-thermal. As can be seen from the energy contents, there is an equipartition in the energy content between the electrons and the ions. This result was obtained from observations by Share & Murphy (1995) and Ramaty et al. (1995) and showed that the energetic ions, rather than playing a minor role in the energy budget as previously supposed, are of comparable importance to the electrons.

A central diagnostic of energetic ions during solar flares are ion abundances that are observed in space. The most complete measurements of elemental abundances in the solar corona come from these observations rather than from photons of the coronal
Figure 1.4: The average abundance enhancement in impulsive solar flares vs. charge-to-mass ratio at $3.2 \cdot 10^6$ K. When averaging over a large number of flares, a pronounced dependence of ion enrichment on the charge-to-mass ratio is observed (Reames, 1998).

plasma. However, impulsive flares turn out not to be a good source for probing coronal abundances. While gradual events accelerate all ions in a fairly equitable manner and therefore retain coronal abundances, impulsive flares exhibit variations in enrichment of the different ion species above $\sim 1$ MeV/nucleon (see reviews by e.g. Lin, 1987; Reames, 1990) and, as recently observed, even below this energy (e.g. Mason et al., 2000, 2002). When averaging over many flares, the abundance enhancements during impulsive flares show a pronounced dependence on the charge-to-mass ratio $Q/A$. The dependence of the major ion components is depicted in Fig. 1.4 for values at $\sim 3 \cdot 10^6$ K (He here refers to the most abundant isotop $^4$He). Specifically $^3$He, which is not shown in Fig. 1.4, exhibits a spectacular enhancement of about $5 \cdot 10^4$ over coronal values. These so-called $^3$He rich events are nowadays identified with impulsive flares and are associated with escaping energetic electrons of about $2 - 100$ keV, which can be observed as type III radio emission. Again, it has to be noted that not each individual event shows this systematic dependence and the plot in Fig. 1.4 is obtained by averaging over dozens of flares. But it is a common feature of all impulsive flares that they exhibit heavier ion abundance enhancement and an enormous $^3$He enrichment.

The coronal values are usually taken from observations during gradual flares, where the abundances of the corona are conserved. Also, a general but not so significant mass ordering is observed. As can be seen from Fig. 1.4 the enhancement of an ion species increases with increasing mass. Recent observations describe rare events where the mass
ordering is violated (Mason et al., 2002). While the latter dependence was suggested to be a result of gravitational settling (Mason et al., 1986), the charge-to-mass dependence is a strong indicator to resonant wave-particle interaction as the responsible accelerating, or at least fractionating, mechanism. However, a dependence as shown in Fig. 1.4 and especially a unique enrichment as for $^3$He suggest that the requirement to reproduce the ion abundances may be the most stringent constraint on an acceleration theory and therefore is of special interest.

1.3 Particle Acceleration in Solar Flares

A number of mechanisms have been proposed as possible acceleration agents for solar energetic particles. They mainly can be divided into three broad classes, namely acceleration by parallel DC electric fields, shocks or stochastic acceleration by plasma waves. The three classes are briefly discussed in the following while the last one is illustrated in more detail for a particular scenario in Section 1.3.4.

1.3.1 Parallel Direct Electric Fields

Perhaps the most direct way to accelerate particles is by large-scale DC electric fields. Most work on this topic has been focused on electrons although ions can also be accelerated by DC electric fields. In an external DC electric field, electrons experience in addition to the electric field force a drag force resulting from Coulomb interaction with the other electrons in the population. The drag force depends on the particle velocity and has a maximum at $v_{te}$ while decreasing with increasing electron velocities above the thermal velocity. The key parameter to describe the balance between these two forces is the so-called Dreicer field and is defined as the field that a thermal electron has to experience in order to gain infinite energy. Two regimes can be distinguished: the applied field is either larger than the Dreicer field $E > E_D$ (super-Dreicer) or smaller $E < E_D$ (sub-Dreicer). In the former case of super-Dreicer fields most electrons of the population will gain energy since the electric field force overcomes the drag force for all particles. Since there is no resulting drag force, the electrons and ions will then be freely accelerated to higher energies. For sub-Dreicer fields there is a critical velocity $v_c$ which separates regimes of limited particle heating ($v < v_c$) and free electron acceleration ($v > v_c$).

A model for electron acceleration by super-Dreicer fields is proposed by Litvinenko (1996, and references therein). Super-Dreicer fields occur in the current sheets of reconnection above magnetic loops. These fields largely exceed the Dreicer fields by typically several orders of magnitude. Whereas the model can yield energies and fluxes consistent with HXR observations, i.e. electrons below $\sim 100$ keV, the problem of replenishment has not yet been fully investigated. Also, it is not very likely that the highest energy electrons can be obtained. Although it is in principle possible to produce 10 MeV electrons by the potential drop assumed, the particles are driven out of the current sheet before they have experienced and gained the available energy. It is therefore not clear
how many electrons will gain the maximum energies.

Sub-Dreicer acceleration has been considered for a long time in the physics of laboratory plasmas. Its application to solar flares was proposed by Holman (1985) and Benka & Holman (1994). There were successes in modeling spectral HXR data by a composite model involving a thermal contribution from the electrons below $v_c$ and a contribution from the runaway population. It turned out that less electrons are needed ($\sim 10^{33} - 10^{34}$ electrons s$^{-1}$) to produce the observed HXR spectra than in a pure non-thermal bremsstrahlung model. However, sub-Dreicer scenarios suffer from a limitation in the electron energy to below about $\leq 500$ keV. The highest observed energies of tens of MeV cannot be readily produced by such a scenario.

Ions can also be accelerated and run away due to electric fields. In the case of $E < E_D$ the ion energy is limited to the range between $0.1 \cdot v_t e$ and $v_t e$. In super-Dreicer fields, ions may be directly accelerated to higher energies. By assuming solar flaring conditions, ion energies of about $\sim 10 - 10^3$ keV can be reached. Again, the production of the highest energy particle is the most severe problem of the model.

Another limitation to both mechanisms is the need to ensure overall neutrality of charge and current in the acceleration region. Since the inferred number of energetic particles requires that a significant fraction of the coronal population, in fact the whole flaring volume, is accelerated, replenishment of the particle source region has to be provided. In addition, the energetic particles correspond to a significant electric current, that has to be neutralized in order not to generate large magnetic fields. While this problem may be less severe for large electric fields ($E \gg E_D$) where accelerated particles can escape from the electric field before a large current can build up, it is acute for the sub-Dreicer case since, due to the weak electric fields, the potential drops have to extend over a wide range in order to gain the observed energies.

A general problem for all kinds of electric field acceleration is the difficulty to reproduce abundance variations as described in Section 1.2.2. Observations require a dependence on the charge-to-mass ratio of the appropriate ion species which cannot be produced by electric field acceleration. Nevertheless, in the case of ion acceleration one can think of an application as pre-accelerator where the fractionating is done by another mechanism later on.

To conclude this section it is pointed out that for both of the mechanisms described above the actual acceleration is only in the direction parallel to the DC electric field and, hence, to the background magnetic field. It is a result of the directed nature of the electric field acceleration, that the particle distribution function will become anisotropic with an excess in the longitudinal direction in the absence of an efficient scattering mechanism.

### 1.3.2 Shock Acceleration

Shock acceleration has been widely invoked to account for energetic particles in solar flares (gradual and impulsive), the solar wind, the earth’s bow shock and termination shock, and at galactic and extra-galactic locations. Cosmic rays are a well known example of an application of shock acceleration (e.g. Zank & Gaisser, 1992). It is now
recognized that there are two quite different mechanisms by which a shock wave can accelerate particles. These are shock drift acceleration and diffusive shock acceleration and are described below.

Shock drift acceleration refers to particles moving along the surface of the shock, gaining energy from the shock’s electric field. Electrons behave approximately adiabatically due to their gyro radius being well below the characteristic shock length scales. There are two main properties for this kind of shock acceleration severely limiting the effectiveness of the process. In the absence of any scattering mechanism, the particles escape along the upstream magnetic field and are lost to the acceleration process. The interaction is therefore reduced to either a transition across the shock front or only one single reflection. In the case of the well known Fermi (see also Section 1.3.3) process of two converging shock fronts, the particle velocities become more and more aligned along the magnetic field as their energy increases with each reflection (Melrose, 1990). This in turn reduces the probability of them being reflected (and further accelerated) rather than transmitted (and lost) on encountering a shock front. Thus the acceleration tends to self limitation in the absence of any scattering mechanism. Second, as was shown by simulations of Wu (1984) and Krauss-Varban et al. (1989), the mechanism is only effective when the angle between the upstream magnetic field and the shock normal is close to 90°. This results in a very limited particle flux (e.g. \(<1\%\) of the electrons). While drift acceleration is very successful at the earth’s bow shock (Wu, 1984; Leroy & Mangeney, 1984) the restrictions in the process are by far too limiting to be considered seriously in solar flares. However, by including a scattering mechanism in the upstream region it was shown by (Decker & Vlahos, 1986) that MeV energies can be achieved, and it is well known that Fermi acceleration is a very efficient mechanism in the presence of a resonant scatterer.

Current ideas on the second type of shock acceleration may be regarded as a modification of the first-order Fermi mechanism described above. A single fast-mode shock can very efficiently accelerate particles in the presence of scattering centers in the upstream and downstream regions. When going to the rest frame of the fluid, on either side of the shock, the scattering centers on the other side of the shock are moving towards the shock front. Hence, every time a particle crosses the shock front, the first encounter with a scattering center is head-on and the energy gain is positive. While being scattered many times across the shock front, acceleration proceeds and the net energy gain is substantial.

While the scattering centers have not been specified, their nature poses a restriction to the diffusive shock acceleration. If wave turbulence is assumed to deliver the necessary scattering, the particles have to fulfill a resonance condition in order to experience efficient scattering. This puts a lower limit on the initial velocities of the various particle species to become accelerated. This so-called ‘injection problem’ makes it difficult to describe particle acceleration in solar flares by diffusive shock acceleration as a single-stage model. Whilst it is less restrictive for ions (Kucharek & Scholer, 1991), in the case of Alfvén turbulence electrons have to exceed a speed of \((m_p/m_e)v_A\) where \(v_A\) is the Alfvén speed. These are already relativistic energies and a pre-accelerator is needed to explain acceleration out of a thermal state. The injection energy is dependent on the
wave turbulence assumed, since it is a result of the resonance condition.

While shock drift acceleration seems to be unlikely as the major accelerator in solar flares, diffusive shock acceleration is a candidate for the acceleration of electrons and ions. The injection problem restricts possible models to two-stage processes that involve a pre-acceleration mechanism, energizing the particles to the required injection energies where they can be caught by resonant scattering at the shock and further accelerated to higher energies. However, the equitable manner in which shocks accelerate different species of ions lacks an explanation for the observed ion abundance anomalies described in Section 1.2.2.

1.3.3 Stochastic Acceleration

The suggestion by Fermi (1949, 1954) that the galactic cosmic rays are accelerated by bouncing off moving magnetized clouds already includes ideas most relevant to stochastic acceleration. They can be summarized in three major ingredients to Fermi’s mechanism. One is that a particle gains energy in a head-on collision with a cloud and loses energy in trailing collisions. The second idea is that if the cosmic rays are moving at random, the probability of head-on collisions is increased due to the relative speed between cloud and particle. The third ingredient is the implicit assumption of isotropy of the particle distribution. A special mechanism is required to maintain an isotropic state since the collisions tend to align the particle velocities along the magnetic field, severely limiting the acceleration.

Fermi-type acceleration, in the sense of processes involving adiabatic heating, also occurs in a variety of other accelerating processes e.g. magnetic pumping (Swann, 1933; Falthammer, 1963), where temporal variations of the magnetic field cause adiabatic heating of trapped particles or transit acceleration (Shen, 1965) by particle diffusion through inhomogeneous magnetic fields. The important features of all these mechanisms are that they involve statistical energy gains and that they require a mechanism that maintains the assumed isotropy of the particles.

While the classic Fermi mechanism involves large-amplitude magnetic compressions interacting in a non-resonant manner with the particles, a richer and more modern approach to stochastic acceleration was introduced by Achterberg (1981) who pointed out that the interaction with low-frequency turbulence (as in the Fermi mechanism) may be described in terms of Čerenkov resonance, whereas the isotropizing scattering by high frequency waves occurs at cyclotron resonances. Fermi-type acceleration can therefore be regarded in these terms as a result of magneto-acoustic turbulence. MHD turbulence consists of a mixture of two modes, both propagating approximately at the Alfvén speed $v_A$: The compressive and isotropically propagating fast-mode, $v_A = \omega/|\mathbf{k}|$ and the torsional parallel propagating Alfvén mode $v_A = \omega/|\mathbf{k}_||$. Many MHD turbulent models involve a magnetic cascade. The cascade is excited on a large characteristic scale length (small $k$). Non-linear interaction transfers the energy dissipation-free through the so-called 'inertial region' to smaller scale lengths (larger values of $k$). The energy finally is dissipated into the particles via the corresponding resonances in the so-called 'dissipation range'. For Fermi-type acceleration only the fast-mode branch is dissipated
via Čerenkov resonance. The Alfvén branch is very inefficiently damped. However, cyclotron resonance with ions has been proposed to dissipate the energy in the Alfvén branch (Miller & Roberts, 1995; Miller & Reames, 1996), resulting in the total energy of the cascade being available for particle acceleration.

The following section presents a realization of such a stochastic scenario. It involves a refined version of the classic Fermi mechanism described in the above and is a successful application of a MHD turbulent acceleration model.

1.3.4 Acceleration by Transit-Time Damping

A broadband spectrum of waves is typically required to stochastically accelerate particles to high energies. Such a spectrum is produced by MHD turbulence that is excited by an outflow with velocity of order of the Alfvén speed from the reconnection site, i.e. the site of primary energy release. The corresponding scale length of the excited waves is given by the size of elementary flux tubes of about $\sim 10^8$ cm (e.g. Petschek, 1964; Priest, 1985; Forbes & Priest, 1987; LaRosa & Moore, 1993). The turbulence is isotropic, consistent with Kolmogorov-like and Kraichnan cascades (Kraichnan, 1965) and transfers the energy to scales where the waves preferentially accelerate electrons out of the thermal distribution and are subsequently damped. The compressive magnetic component of the cascade, i.e. the fast-mode branch, interacts with the electrons through Čerenkov resonance. This process is called transit-time damping (Fisk, 1976; Stix, 1992) and is the magnetic equivalent of Landau damping (which involves Čerenkov resonance with parallel electric fields). This mechanism was also termed small-amplitude Fermi acceleration (Achterberg, 1981) in the literature. The wave-particle interaction in transit-time damping is between the particle's magnetic moment and the magnetic field gradient of the wave. As in Fermi acceleration the particles are reflected by the magnetic compression. But, contrary to the classic Fermi mechanism, only a few particles are reflected by the small amplitudes of the waves. In the limit of very small amplitudes, the parallel particle speed in the wave frame moving along $B_0$ with the parallel phase speed must be close to zero for reflection to occur. This is just the resonance condition for Čerenkov resonance $k_\parallel v_\parallel = \omega$. The interaction is therefore of a resonant nature as opposed to the non-resonant mirroring in the case of the classic Fermi mechanism.

An essential point is the necessity of a spectrum of waves for efficient acceleration. Interaction with one wave results in an energy change of the particle. Dependent on the velocity difference, the particle gains or loses energy. The change in $v_\parallel$ removes the electron from resonance with this wave and in the case of only a single small-amplitude wave the process stops. In the case of a wave spectrum, the change in $v_\parallel$ brings the electron into resonance with a neighboring wave, resulting in another energy gain or loss, and so on. The particle therefore 'jumps' from resonance to resonance through the spectrum and can systematically gain energy (Karimabadi et al., 1992). As in the Fermi mechanism, head-on encounters are more probable than trailing encounters and a net energy gain is achieved. The isotropy in the propagation angles of the fast-mode waves, which is a result of the turbulent cascade, ensures a wide span of available parallel phase velocities $v_\parallel = v_A / \cos \theta$ of the waves, where $\theta$ denotes the propagation angle with
respect to the background magnetic field.

Another important aspect in transit-time damping is the need for isotropization, since Čerenkov resonance is only capable of transferring energy in the parallel direction. As already mentioned in Section 1.3.3, isotropization is a general requirement for effective acceleration in Fermi-type scenarios. The absence of a pitch-angle scatterer leads to a decrease in acceleration efficiency due to a systematic reduction of pitch-angles. The faster the particles become, the more they are aligned along the magnetic field and higher parallel phase speeds of the waves are required for resonant interaction (i.e. waves at propagation angles close to 90° with respect to \(B_0\)). This reduces the available wave density for acceleration and the process ceases. Pitch-angle scattering and resulting isotropization therefore greatly increases the portion of the wave spectrum that can be sampled by the electrons and enables large energy gains.

Since the thermal speeds of all ion species lie well below the Alfvén velocity, ions cannot be directly accelerated out of the thermal distribution by the mechanism described above. In this scenario, heavy ions are accelerated via cyclotron resonance by the Alfvén waves in the torsional branch of the cascade. As the waves cascade from lower to higher frequencies, they will encounter Fe ions first and strongly accelerate them. Due to the low Fe abundance, the waves are not significantly damped. They further cascade to higher frequencies, accelerating ions of the Ne group, until the rest of the total Alfvén wave energy is consumed by the ions at the O group resonances. The heavy ion abundance enhancements described in Section 1.2.2 are therefore produced by the layering of the cyclotron resonances of the various ions and the loss of energy in the cascade (Miller, 1998).

Proton acceleration poses a special problem since the total energy of the Alfvén branch is consumed by the ions with resonances at lower frequencies: The proton gyro-frequency lies above the \(^4\)He gyrofrequency where the cascade stops and no power is left in the waves. In the original model, mode conversion of waves in the fast-mode branch has been proposed to account for proton acceleration. It is not yet clear how efficient this mode conversion is and if it is capable of producing protons up to the energies observed. However, in Chapter 3 it is proposed that the EFI may be responsible for transferring energy from the fast-mode branch via the electrons to the protons. This may result in at least a pre-heating of the protons that will allow them to be caught by resonance either with Alfvén waves at lower frequencies or by fast-mode waves via transit-time damping, analogously to the electrons.

The unique enhancement of \(^3\)He cannot be explained as a result of MHD turbulence and transit-time damping or cyclotron resonance with Alfvén waves. It requires a completely detached mechanism in this model. A possible acceleration mechanism for \(^3\)He was proposed by Temerin & Roth (1992) and adapted by Miller & Viñas (1993): Electron beams of about 5 keV excite so-called \(H^+\) EMIC waves in a narrow range around the \(^3\)He resonance frequency. No other resonance lies in the frequency range of the waves and \(^3\)He is therefore selectively accelerated. A severe problem of this mechanism is on one hand the existence of such beams and on the other hand, that such a beam would be eroded by a bump-in-tail instability on time scales much shorter than the growth of the \(H^+\) EMIC waves (~ 4 orders of magnitude). A closer description of this model and
an alternative mechanism to accelerate $^3$He is described in Chapter 5.

The model for acceleration by transit-time damping fulfills most of the requirements for a successful flare particle accelerator. Not only can it describe electron and ion acceleration as a direct consequence of an MHD turbulent cascade, it also predicts ion abundance enhancements that are observed in space. However, isotropization of the electron population, the mechanisms for proton acceleration and $^3$He enrichment are not satisfactorily explained in the model. All three issues are addressed in this work and are outlined in the following section.

1.4 Outline

Some unexplained elements of flare particle acceleration are addressed in this work. On one hand, observations were carried out to find further evidence for the location of the accelerator, while on the other hand, several parts of the acceleration process have been modeled analytically and in computer simulations.

First, observations are presented that were carried out to locate the acceleration region of escaping energetic electrons and to confirm the canonical flare picture sketched in Section 1.1. This was done in collaboration with the Radioheliograph Group in Meudon, Paris. Simultaneous observations of solar type III radio emission associated with metric spikes were carried out with the Phoenix-2 spectrometer in Zurich and the Nançay Radioheliograph (NRH). An attempt was made to follow and extrapolate the type III emission caused by escaping energetic electrons to lower altitudes and to compare it with the source region of radio spike emission. The goal was to answer the question of whether metric spike emission can come from the primary acceleration region, and if so is it above the loop as expected from HXR observations (Masuda et al., 1994)? The results of this investigation are presented in Chapter 2.

The theoretical investigations start with the analysis of the electron distribution function during solar flares. The main common feature of the accelerated population is found to be an anisotropy in velocity space with $T_\parallel < T_\perp$. Except diffusive shock acceleration, all of the above mentioned acceleration mechanisms directly yield such an anisotropy. It is established that such a distribution function becomes unstable to the so-called Electron Firehose instability (EFI) under solar conditions. Chapter 3 presents an instability analysis in linearized kinetic theory and describes the properties of the excited waves. In addition, a new mode has been discovered propagating at oblique angles with respect to the background magnetic field.

The resulting anisotropy in many acceleration mechanisms destroys the efficiency of the accelerator. While, in many of the aforementioned accelerators, an external source for pitch-angle scattering of the electrons (and hence isotropization of the electron distribution) is proposed, the EFI is a potential scattering agent as an inherent ingredient of electron acceleration. The great advantage in comparison to a scatterer from an external source lies in the fact that the EFI is an intrinsic feature of the acceleration process itself. The abilities of the EFI to pitch-angle scatter or better, transfer energy to higher perpendicular velocities, is investigated by using an analytical approach and
Figure 1.5: This flow chart shows the theoretical subjects of particle acceleration during impulsive solar flares that have been addressed in this work (grey shaded). They are embedded in the exemplary model of acceleration by transit-time damping, which could be replaced by any of the aforementioned accelerators. The observational part covers an aspect of accelerated electrons and is not indicated in the chart.

test particle simulations. The results of this work are presented in Chapter 4.

Another application of the EFI to solar flare energetic particles is found in its ability to accelerate and enhance $^3$He. Again, as a necessary consequence of accelerated electrons, Electron Firehose (EF) waves grow and influence the flaring plasma. The excited EF spectrum is very stable for a variety of plasma parameters that could arise during solar flares. These spectral properties allow $^3$He to be accelerated preferentially with respect to $^4$He to the observed energies. Analytical calculations and test particle simulations were applied to the problem and the results of these are presented in Chapter 5. For the first time it is therefore possible to describe $^3$He acceleration and enrichment as a direct result of the bulk electron energization, also explaining the strong correlation between electron and $^3$He rich events.

The work is concluded by summarizing the findings and giving an outlook on future projects resulting from this thesis in Chapter 6.
The flow chart of the acceleration scenario in Fig. 1.5 summarizes the work presented here. It shows the embedding of the results in the model of acceleration of electrons by transit-time damping. This model was chosen as a representative scenario, although it has to be noted that it can be replaced by most of the acceleration mechanisms described above. In this view, the EFI provides two inherent ingredients: on the one hand it enables efficient acceleration by isotropizing electrons and on the other hand it leads to accelerated and enriched $^3$He, both as intrinsic results of electron bulk energization in impulsive solar flares.
Chapter 2

Spatial Analysis of Solar Type III Events Associated with Narrow Band Spikes at Metric Wavelengths

ABSTRACT: The spatial association of narrow band metric radio spikes with type III bursts is analyzed. The analysis addresses the question of a possible causal relation between the spike emission and the acceleration of the energetic electrons causing the type III burst. The spikes are identified by the Phoenix-2 spectrometer (ETH Zurich) from survey solar observations in the frequency range from 220 MHz to 530 MHz. Simultaneous spatial information was provided by the Nançay Radioheliograph (NRH) at several frequencies. Five events were selected showing spikes at one or two and type III bursts at two or more Nançay frequencies. The 3-dimensional geometry of the single events has been reconstructed by applying different coronal density models. As a working hypothesis it is assumed that emission at the plasma frequency or its harmonic is the responsible radiation process for the spikes as well as for the type III bursts. It has been found that the spike source location is consistent with the backward extrapolation of the trajectory of the type III bursts, tracing a magnetic field line. In one of the analyzed events, type III bursts with two different trajectories originating from the same spike source could be identified. These findings support the hypothesis that narrow band metric spikes are closely related to the acceleration region.

2.1 Introduction

Milliseccond narrow band radio spikes are structures in the radio spectrum of the Sun forming a distinct class of flare emission. The term 'narrow band, millisecond spikes'
refers to short (few tens of ms) and narrow band (few percent of the center frequency) peaks in the radio spectrogram. They can be observed in the range of 0.3 to 8 GHz and occur mainly during the impulsive phase of a solar flare. Since the spike emission is often associated with enhanced hard X-ray emission (Benz & Kane, 1986; Güdel et al., 1991; Aschwanden & Güdel, 1992) it is suspected that spikes are closely related to the actual process of energy release in solar flares.

The short duration and the narrow bandwidth suggest a small source size and therefore a high brightness temperature (up to $10^{15}$ K). Only a coherent mechanism can account for the emission but none of the proposed mechanisms is generally accepted. Proposed emission mechanisms are the electron cyclotron maser (Holman et al., 1980; Melrose & Dulk, 1982; Aschwanden, 1990; Robinson, 1991) and upper hybrid or z-mode instabilities combined with wave-wave coupling (e.g. Zheleznyakov & Zaitsev, 1975; Vlahos et al., 1983; Tajima et al., 1990; Güdel & Wentzel, 1993).

A subclass of spikes originally found at metric wavelengths correlates with type III bursts (Benz et al., 1982). They occur in clusters usually at frequencies slightly higher than the start frequency of the type III burst and are located in a dynamic spectrogram at the intersection of the extrapolated type III and the spike frequency. Although they may be slightly shifted in time (in a positive or negative direction) they significantly correlate with the extrapolated type III burst (Benz et al., 1996). Since the spike location in the spectrogram is close to the extrapolated type III burst it is a reasonable assumption that spike and type III radiations are emitted at the same characteristic frequency.

This type of radio emission has been called 'metric spikes' in the literature (e.g. Güdel & Zlobec, 1991) and it is still unclear whether the spikes in the metric and the decimetric range belong to the same class of events.

Previously published spatially resolved observations of metric spike events (Krucker et al., 1995, 1997) found the spike sources at high altitudes and suggest a model of energy release taking place in or close to the spike sources. Escaping beams of electrons cause the type III emission. Thus a scenario is conceivable in which the spikes may be a direct signature of the accelerator.

Using two-dimensional, spatially resolved data from the Nançay Radioheliograph (NRH), it is possible to reconstruct the spatial configuration of the event and the relative position of the spike source with respect to the type III trajectory. The main purpose of this work is to test whether the geometry of the events supports the picture mentioned above.

### 2.2 Instruments

#### 2.2.1 The Phoenix-2 Spectrometer

Since 1998 the Phoenix-2 spectrometer operated by the ETH Zurich has been continuously recording radio data from sunrise to sunset. The intensity and polarization are digitally measured by the frequency-agile receiver in the range from 0.1-4.0 GHz at a
time resolution of 500 μs for single channel measurements. The receiver consists of 4000 channels at 1, 3 or 10 MHz bandwidth, from which a reduced number can be freely selected. A full description of the instrument can be found in Messmer et al. (1999).

The data used herein were recorded at a bandwidth of 1 MHz and a time resolution of 100 ms in the frequency range from 220 to 550 MHz, chosen in collaboration with the Nançay Radioheliograph.

### 2.2.2 The Nançay Radioheliograph (NRH)

In July 1996 the Meudon Observatory began daily observations with the improved 2D imaging radioheliograph in Nançay (France). The instrument is described by Kerdraon & Delouis (1996). Five frequencies in the range from 150-450 MHz can be observed simultaneously at a maximum number of 200 images per second. The antennas are organized in two perpendicular arrays and digitally correlated on 576 channels resulting in measurements of Stokes I and V. The observing bandwidth is 700 kHz.

Data were taken with a frequency configuration of 164.0, 236.6, 327.0, 410.5, 432.0 MHz at a mean time resolution of 125 ms.

For older data the NRH provides one-dimensional scans of the corona at five frequencies, with both its east-west and north-south branches (The Radioheliograph Group, 1993). The integration time was 1 s. The NRH then observed at 164, 236.6, 327, 408 and 435 MHz.

### 2.3 Observation and Method

Data from observations of Phoenix-2 were used to identify type III events associated with metric spike emission. The choice of events was restricted by the observing frequencies of Nançay. At least one frequency was required to locate the spike emission, and the type III emission was required to be visible at least in two frequencies. The following four events were found in the recent data:

- **99/09/27:** In the time range from 09:53.8 to 09:57.2 UT a type III burst group occurred and was accompanied by a cluster of metric spikes in the frequency range 300 to 370 MHz.

- **00/05/29:** Type III burst group in the range from 220 - 420 MHz and metric spike emission from 390 to 430 MHz in the time range from 12:25.00 to 12:29.80 UT.

- **00/06/21:** Metric spikes in the frequency range from 340 MHz to 450 MHz, accompanied by a type III burst group from 220 MHz to 450 MHz. The event time is 10:12.4-10:12.5 UT.

- **00/06/22:** Type III burst group in the range from 220 MHz to 510 MHz, whereas the metric spikes are located between 400-420 MHz at the beginning of the event and
between 430-500 MHz later on. The whole event lasts from 13:41.5 to 13:44.5 UT.

In addition, three separate events on 92/08/18, already published in Krucker et al. (1997) using only VLA observations at 333 MHz have been investigated using data from the previous heliograph in Nançay. The events show metric spikes in the frequency range from 310 MHz to 360 MHz and type III activity from around 300 MHz down to below 40 MHz. The time ranges for the events are 13:43:50 to 13:44:10 UT, 14:01:47 to 14:02:00 UT and 14:14:35 to 14:14:40 UT, respectively.

2.3.1 Phoenix-2 Observations

The dynamic spectrum from 220 MHz to 530 MHz of the analyzed part of the 00/06/22 event is shown in Fig. 2.1 (top). In the first half of the the event (13:41:40 to about 13:42:08 UT) the metric spike emission is mainly located in the range from 370 MHz to 470 MHz. At lower frequencies the drifting structures of type III emission can be seen. In the second half, the spike emission is shifted to higher frequencies (430 MHz to 510 MHz).

In the middle panel of Fig. 2.1 a selected light curve at 432.0 MHz observed by the Phoenix-2 spectrometer is displayed. The data were recorded with a sampling period of 100 ms and an integration time of 0.5 ms. According to Güdel & Benz (1990) the expected mean time duration of a single spike event at 432.0 MHz is 0.062 ± 0.004 s. Therefore the single spike events are not resolved in time. Following Messmer & Benz (1999), the minimum bandwidth of spikes is given by \( \sim 0.4\% \) of the center frequency yielding a value of 1.73 MHz at 432.0 MHz. Although this value is close to the spectral resolution of 1 MHz of Phoenix-2, the spikes are resolved in frequency and can be well identified.

No polarization information is available from Phoenix-2 data for this event nor for the 00/06/21 and the 00/05/29 events since the spectrometer was run with a linear feed during this time.

The description of the observation mode of the old Phoenix spectrometer for the 92/08/18 event can be found in Krucker et al. (1997)).

2.3.2 Nançay Observations

The bottom panel in Fig. 2.1 shows the light curve at 432.0 MHz of the same event observed by NRH. The data were recorded with an integration time of 125 ms at a sampling rate of 125 ms. Since the integration time of NRH is longer than of Phoenix-2, the noise level of the dataset is lower.

The software used to analyze the 2D NRH data offers a source tracking routine that allows the user to determine location and size of a radio source in the 2D image of the Sun. The relevant source was identified from Phoenix-2 observations by simultaneously determining the flux in the NRH data and comparing it to the full sun observation of
Figure 2.1: Top: Spectrogram observed by Phoenix-2 on 00/06/22. White regions correspond to enhanced flux; the frequency axis is from top to bottom. Middle: Light curve of metric spikes recorded by the Phoenix-2 spectrometer at a single frequency (432.0 MHz). The arrow in the top panel indicates the position of 432.0 MHz. Bottom: NRH light curve of metric spikes at the same frequency.
Phoenix-2 (cf. Fig. 2.1 middle and bottom). The size of the source is determined by fitting an ellipse with minimal perimeter at half height of the peak flux.

The error of the source centroid consists of two parts: A discretization error and a statistical error. The first one stems from the NRH data being gridded. The centroid position determined by the NRH software lies on a grid and the distance between two grid points is \( \sim 0.016R_\odot \). Hence, there is a minimal error which is given by half the diagonal of a pixel \( \sim 0.011R_\odot \).

The statistical uncertainty of the centroid position is given by the observed radius, \( \Phi \), of the source times the noise to flux ratio. This ratio is constant according to the radiometer equation, and the resulting error \( \Delta \) can be written as

\[
\Delta = \frac{\sigma_{\text{max}}}{F_{\text{max}}} \Phi = \frac{\sigma}{F_{\text{bg}}} \Phi
\]  

where the index \( \text{max} \) refers to the peak values, \( \sigma \) is the noise level at the background and \( F_{\text{bg}} \) is the background flux. The noise can be determined from the background level at the source position before or after the event. In most cases, the statistical error was smaller than the discretization error, thus the latter one was used. Otherwise Eq. (2.1) yields the error bar.

In the case of the 92/08/18 event, the integration time in the NRH data is so long that the accuracy of the centroid determination is given by the pixel size. This has been verified in the following way: 1D scans were made at 164 and 327 MHz during five observing runs on Cyg A between 6 and 31 August 1992 and analyzed in the same way as the solar data. At 164 and 327 MHz, Cyg A appears as a simple source when observed with both the east-west and north-south array. The measured centroid positions contain a random part and a slowly varying offset, which was evaluated by a polynomial fit. The analysis yields a statistical uncertainty of \( \sim 1 \) pixel (error bar) for both frequency channels and a systematic offset of one pixel in the interval of hour angles corresponding to the solar observations on 92/08/18.

### 2.3.3 Association of Spikes and Type III Bursts

The identification of a single type III burst was mainly made by visual identification in the spectrogram. Only the stronger type III bursts of a group have been chosen for the analysis in order to ensure that they can be continuously traced back to the start frequency. If doubt arises, the correlation function of two light curves at the relevant frequencies was computed and the burst was identified by correlation peaks. In case of more than one correlation peak, consistency with the drift rate between other frequencies was the criterion. However, it is still possible that two type III bursts with only a slight difference in drift rate crossover in the spectrogram and the associated sources do not belong to the same burst. As described in Sect. 2.4 the sources at each frequency of the analyzed events are very stable. Either each of the bursts in the group has a crossover with another type III on the “wrong” field line, which is not very probable, or the crossing type III burst occurs on the same field line and the source location is de facto not distinguishable from the source of the original burst.
By extrapolating the type III burst to the spike frequency, the associated spikes were identified. However, since the spikes can be slightly shifted in time with respect to the extrapolated type III burst (Benz et al., 1996), a time window of about ±0.2 s has to be used in order to take all possible associated spikes into account. The positions of all spikes inside the time window were included in the analysis.

Simultaneous emissions at two frequencies have been admitted if two of the observed frequencies lie in the spike range.

2.4 Results

\[ \text{Figure 2.2: 3-dimensional view of one of the bursts in the 99/09/27 event. The projection on } A \text{ corresponds to the upper left panel in the plots of Figs. 2.3 to 2.6, 2.8 and 2.10, the projection on } B \text{ to the upper right and the projection on } C \text{ to the lower right panel.} \]

The results for the four recent events (99/09/27, 00/05/29, 00/06/21, 00/06/22) are presented in chronological order whereas the 92/08/18 events will be analyzed separately. In order to display three dimensional situations in two dimensions, height projections on the solar equatorial plane and the meridian plane which we define by the earth and the solar poles have been chosen for representation in Figs. 2.3 to 2.8 and 2.10 (see Fig. 2.2). The upper right panel of each plot depicts the actual observations of Nançay as described in 2.4.1 whereas the upper left and the lower right panel result from the 3-dimensional reconstruction described in Sect. 2.4.3. The symbols representing the observing frequencies are chosen as follows: \( * = 432.0 \text{ MHz}, + = 410.5 \text{ MHz}, \times = 327.0 \text{ MHz}, \square = 236.6 \text{ MHz} \) and \( \Delta = 164.0 \text{ MHz} \).
The lines connecting the source centroids display how the sources are related to each other in the spectrogram. For reasons of visualization the observed source centroid positions in the 92/08/18 events have been interpolated by 3-dimensional splines to smooth curves.

The sources do not vary significantly in the course of the event for all analyzed events. Hence, for every event, one single representative burst out of the group has been plotted, although the analysis contains many more identified type III bursts.

### 2.4.1 Observation

References to Figs. 2.3 to 2.8 and 2.10 in this section always refer to the upper right panel displaying the square shown in the overview (lower left panel).

**99/09/27:** Five type III bursts were identified at 164.0 MHz during this event. Two of the NRH frequencies (164.0 MHz, 236.6 MHz) lie in the range of the type III bursts and one (327.0 MHz) in the spike range. The event occurred on the first quadrant of the solar disc. The source positions at each frequency vary only within the error bars, indicating a stable magnetic configuration during the event. Therefore it can be assumed that the bursts occurred practically on the same magnetic field lines. The bursts of this group are depicted in Fig. 2.3 (upper left plots).

**00/05/29:** Three type III bursts were identified at 164.0 MHz in this group occurring between 12:28:46 and 12:28:49 UT. All five observing frequencies of NRH could be used for this event. Three frequencies (164.0 MHz, 236.6 MHz, 327.0 MHz) lie in the range of the type III emission and two (410.5 MHz, 432.0 MHz) in the spike range. The spatial variation of the sources among the different bursts is within the error bars and therefore the three bursts are assumed to have occurred approximately on the same magnetic field line. The upper right plots in Fig. 2.4 display the burst out of this group.

**00/06/21:** Between 10:12:28 and 10:12:44 UT six type III bursts were identified. Again all five frequencies of NRH could be used for the analysis. Three frequencies (164.0 MHz, 236.6 MHz, 327.0 MHz) lie in the range of the type III emission and two (410.5 MHz, 432.0 MHz) in the spike range. The event occurred close to the solar limb in the first quadrant. The sources at 164.0, 236.6, 327.0 and 410.5 MHz are stable within the error bars. The source at 432.0 MHz exhibits a larger spatial variation than observed in all other events. The type III source at 327.0 MHz is the most difficult one in the whole dataset to fit with a field line connecting the sources at all frequencies. The plots in Fig. 2.5 display the six bursts of this group.

**00/06/22:** Five bursts out of this group have been analyzed. 236.6 MHz and 327.0 MHz lie in the type III frequency range and 410.5 MHz and 432.0 MHz lie in the spike range. The spatial variation of the sources is negligible, and as in the other events the extrapolated type III trajectory intersects the observed spike source. These bursts of this event are shown in Fig. 2.6.
Figure 2.3: All five well identified type III bursts of the 99/09/27 event. This plot and the following ones are built as follows: Lower left: global position of the event on the Sun. The small square indicates the image size presented in the upper right quadrant, where the positions observed by the NRH are given. The square in the upper right corner represents the grid size of the NRH image. The positional error of the source centroids are indicated by ellipses. Upper left: Projection of the sources on the meridian plane (view seen from an observer West of the sources). Lower right The sources in the projection on the equatorial plane, showing the view of an observer North of the Sun (cf. Fig. 2.2).
Figure 2.4: The three well identified type III bursts of the 00/05/29 event. For a description of the plot see caption of Fig. 2.3.
Figure 2.5: The six identified type III bursts in the 00/06/21 event. The plots are described in the caption of Fig. 2.3.
Figure 2.6: The five well identified type III bursts in the 00/06/22 event. The plots are described in the caption of Fig. 2.3.
Three events of type III bursts with associated metric spikes were identified on 92/08/18 starting at 13:44 UT, 14:02 UT and 14:14 UT, respectively.

13:44 UT: During the first event, three type III bursts were identified at 164.0 MHz. Two more NRH frequencies could be used for the analysis: 236.6 MHz in the type III range and 327.0 MHz in the spike range.

The scattering of the source centroid positions during the event lies within the error bars for each frequency. The bursts of this group are displayed in Fig. 2.7.

14:02 UT: This event is of special interest since it exhibits simultaneous sources at 164.0 MHz at different locations. Only one type III burst could be clearly identified in the spectrogram. At 164.0 MHz two sources were found: a stronger one more to the East and a weaker one more to the West. At 236.6 MHz only one source position

Figure 2.7: All three bursts from the first type III group at 13:44 UT, 92/08/18. A description of the plots is found in the caption of Fig. 2.3.
Figure 2.8: Burst in the second event at 14:02 UT, 92/08/18. A description of the plots is found in the caption of Fig. 2.3.

Figure 2.9: Light curve of the 92/08/18 14:14 UT event at 164.0 MHz. The Roman numerals refer to the four identified type III bursts.
Figure 2.10: Left: Reconstructed trajectories of bursts I and II of the 92/08/18 14:14 UT event. The Roman numerals refer to the first two bursts identified in Fig. 2.9. Right: Burst III and IV of the 92/08/18, 14:14 UT event. In both plots the trajectories have been 3-dimensionally spline interpolated in order to make the distinction between error bars and trajectory easier. The symbols are the same as described in the caption of Fig. 2.2.

could be clearly identified. This position is consistent with only one of the positions at 164.0 MHz. Due to the possible close positions of the sources and the time resolution of 1 s, a weaker second source at 236.6 MHz would be difficult to detect.

At 327.0 MHz, in the spike range, one source was found. Hence, in Fig. 2.8, both 164.0 MHz sources have been plotted whereas only one source at 236.6 MHz is displayed. The more probable trajectory in this configuration is depicted. Nevertheless, since the disconnected 164.0 MHz source is the stronger one and is without doubt part of the event, the situation is interpreted as two type III bursts propagating on different field lines, giving evidence of electrons being injected into different coronal structures from one single spike source.

14:14 UT: The light curve at 164.0 MHz for this type III group is depicted in Fig. 2.9. Four bursts have been identified and labeled I-IV. Special attention has been given to this event for the following reason: Bursts I and IV occurred at a location significantly different from the position of bursts II and III. By analyzing the situation at 236.6 MHz, two sources were found with less, but still significant, spatial separation than at 164.0 MHz. At 327.0 MHz, a frequency lying in the spike range, the positions of the associated spike sources coincide.

The superposition of the radio centroid positions on the Yohkoh-SXT image (Figure 2.10) suggests that the two electron beam trajectories inferred from the radio data correspond to magnetic field lines with different connectivities. The eastward oriented trajectory projects above loops that connect the trailing part of the active region (NOAA 7260) in the West with the leading part of an active region in the eastern hemisphere.
(NOAA 7264). The compact northward oriented trajectory is consistent with electron beams being guided by large-scale magnetic structures bending westward, possibly toward the leading part of AR 7260. The trailing part of AR 7260 had a complex magnetic polarity (δ spot) which produced several flares in August 1992 (cf. Leka et al., 1996) and is thus a plausible site for electron acceleration. Although the geometry is quite similar to the event at 14:02 UT, it exhibits a different situation since the bursts are consecutive and not simultaneous. Therefore, the identification and association of the sources was easier, and at every frequency the single bursts could be well identified.

![Figure 2.11: Yohkoh SXT picture overlaid with radio sources. The SXT image was taken with the Al.1 filter at 16:29:07 UT. The integration time was 2.668 s. Two of the 92/08/18 event radio bursts are displayed (see also Fig. 2.10, right panel). The symbols are the same as described in the caption of Fig. 2.3](image)

2.4.2 Polarization

With the exception of the limb event on 00/06/22, significant circular polarization was detected in all cases. The degree of polarization was moderate (~ 10%) for the type III emission, and much stronger (up to 90-100%) for the spikes. The sense of the polarization was the same for spikes and type III bursts, consistent with the results found by Benz et al. (1996) and references therein.
2.4.3 Spatial Reconstruction

Coronal Density Models

Assuming emission either at the fundamental or harmonic of the plasma frequency, the height of the radio source can be determined from a coronal density model $n_e(h)$ via

$$\omega \approx a \cdot \omega_p = \left( \frac{4\pi e^2 n_e(h)}{m_e} \right)^{1/2},$$

where $h$ is the height above the solar photosphere, and $a$ is about one or two for fundamental or harmonic emission, respectively.

Three different atmospheric models have been used and compared: An exponential atmosphere, the $10^\times$Baumbach-Allen model (Baumbach, 1937; Allen, 1947) and an atmosphere in hydrostatic equilibrium with thermal conductivity (Lang, 1980). The latter two are given by

$$n_e^{BA}(h) = 10 \cdot 1.55 \cdot 10^8 \left( \frac{h}{R_\odot} \right)^{-6} \left[ 1 + 1.93 \left( \frac{h}{R_\odot} \right)^{-10} \right] \text{ cm}^{-3} \quad (2.3)$$

for the Baumbach-Allen model and

$$n_e^{TC}(h) = n_0 \left( \frac{h}{h_0} \right)^{-2/7} \exp \left[ -\frac{7 h_0}{5 H_n} \left\{ 1 - \left( \frac{h}{h_0} \right)^{-5/7} \right\} \right] \quad (2.4)$$

for the thermally conducting corona. $n_0$ and $h_0$ are reference values for density and according height, $H_n$ is the scale height, corresponding to the scale height of the exponential density model. The reference values were chosen according to observations of Trottet et al. (1982) and Suzuki & Dulk (1985). Sources at 160 MHz have been found at heights of about 0.5 solar radii above the photosphere which is reproduced by the model under the assumption of harmonic emission. The resulting density from all three models is depicted in Fig. 2.12.

Although the solar corona is not expected to have a spherically symmetric density distribution, the models are considered to apply to single magnetic flux tubes which are supposed to be the type III guiding structures in the corona. Being anchored in the parent active regions, the flux tubes can exhibit curvature and significantly deviate from radiality. Nevertheless, the density within an individual flux tube can be assumed to follow a model such as described above.

The 3-dimensional position of the radio sources is assumed to be at the intersection of the line of sight and the sphere defined by the density model (Equation 2.2). As all the sources were observed on the disc, there was no ambiguity. Propagation effects and their impact on the source positions are discussed in Sect. 2.5.

If not mentioned otherwise, the exponential density model and $a = 2$ has been used in the following. It is still controversial whether structureless type III bursts (i.e. no fundamental–harmonic pairs) are emitted at the fundamental or the harmonic of the
Chapter 2. Type III Bursts and Metric Radio Spikes

Figure 2.12: Coronal density vs. distance from the Sun center. Three different models have been plotted. The scale height and reference values of density and height for the exponential and thermal conductivity model have been chosen to be $H_n = 7.5 \cdot 10^9 \text{cm}$, $n_0 = 3.36 \cdot 10^8 \text{cm}^{-3}$ and $h_0 = 2.21 \cdot 10^{10} \text{cm}$. The horizontal dotted lines indicate the densities corresponding to the observing frequencies of Nançay assuming emission at the harmonic of the plasma frequency.

plasma frequency. This question shall not be discussed here and we refer the reader to the review of Suzuki & Dulk (1985) and references therein.

The results of the 3-dimensional reconstruction are shown in the upper left (side-view) and lower right (top-view) panels of each plot in Figs. 2.3 to 2.8 and 2.10. The component towards Earth is the projection of the radio source’s height on the axis perpendicular to the plane of the sky (cf. Fig. 2.2).

It is the most important result that in all cases analyzed here, the observed locations of the spikes coincide in a smooth and natural way with the expected position of radio emission at the corresponding frequencies from extrapolating the type III trajectory to lower altitudes.

2.5 Discussion

In the following we discuss the findings in the previous section addressing the 3D reconstruction of the bursts, including possible influence of radio wave propagation effects, and the interpretation of spikes being a signature of the accelerator.
2.5.1 Source Locations

It is obvious that the reconstructed source locations depend on the chosen coronal density model in terms of absolute heights. The same is the case for the choice of emission at the harmonic rather than at the fundamental of the plasma frequency. However, the relative positions which are of major interest in this work are not altered by changing either the atmospheric models nor the characteristic emission frequency. The trajectories may be stretched and shifted in height but the topology of the burst remains the same.

The difference in polarization degree between type III (moderate) and spikes (strong) may indicate another emission frequency for the spikes than occurs for the type III burst (e.g. fundamental for spikes and harmonic for type III). This would shift the spike source to lower altitudes with respect to the type III burst without changing the burst topology. The analysis of Benz et al. (1996) shows that the spikes correlate with the extrapolated type III burst to higher frequencies (i.e. the spike frequency). In case of significantly different emission frequencies for spikes and type III burst, a systematic time offset would be expected.

A common feature of all analyzed events (Figures 2.3 to 2.8 and 2.10) is a bending in East-West direction of the trajectories towards the line of sight to the observer, independent of the location on the solar disc. In North-South direction the data also exhibit non-radiality, but no general trend of deviation was observed. Besides this result being an indication for curved magnetic field structures, there are propagation effects that could produce an apparent bending of the magnetic field lines:

- Refraction in the corona can shift the apparent source location.
- The radio emission may be scattered in the corona during propagation to the observer and therefore can produce an apparent image of the source at a location different from the true position (Arzner & Magun, 1999, and references therein).
- Type III emission may be beamed along the trajectory of the electrons, selecting only those type III bursts propagating towards the observer within the beaming cone (Caroubalos et al., 1974; Caroubalos & Steinberg, 1974).

All events have been analyzed by overlaying EIT images, as shown in Fig. 2.11 for the 92/08/18 event with an SXT image. Comparable indications of magnetic field structures have been found that can explain at least part of the shift of the projected type III positions towards the observer by the guiding magnetic field. E.g., on 00/05/29, the only other event near central meridian, the observed configuration is consistent with sources in an open flux tube that is anchored in the leading part of the active region as seen by EIT and is part of the field lines which project northward onto the disk. Projection effects make association with active regions difficult in the event of 00/06/20. Nevertheless, the radio sources on 00/06/21 have virtually the same north-south coordinate at all frequencies, and the more they are shifted to the east, the lower the frequency. This is expected for an east-westward extending flux tube anchored in
the leading part of the underlying active region, and this interpretation is consistent with the EIT image.

We thus believe that the observations lend plausibility to the assumption that the radio source positions are mainly affected by the magnetic field structure and to a lesser extent by propagation effects. This conclusion is also supported by the work of Pick & van den Oord (1990) and references therein.

### 2.5.2 Spike Location and Acceleration

![Figure 2.13: Sketch of possible locations of the acceleration region (I, II, III) with respect to a type III burst (labeled A) and an associated spike source (asterisk). A second type III (labeled B) is displayed in case of two simultaneous bursts.](image)

To inject electrons on the field line guiding the type III burst, the acceleration region must be located close to or on the field line itself. Figure 2.13 displays a sketch depicting possible locations of acceleration with respect to a type III burst (A) and a spike source. A priori, there are three different positions for the acceleration region consistent with the present observations: it lies between the type III burst and the spike source (case I), below the spike source (case III) or coincides with the spike source (case II). Assuming that the spike emission is caused by the same acceleration event, location I can be excluded by analyzing spectral radio observations. For the acceleration region to lie in
between the type III and the spike source, the time of the actual acceleration must lie within the time interval defined by the intersection of the extrapolated type III with the spike frequency and the start of the type III emission. A systematic time delay of the spike source with respect to the acceleration event must be observed, caused by the travel distance of the electrons generating the spike emission. According to Benz et al. (1996) the spikes correlate with the intersection of the extrapolated type III and the spike frequency itself and no systematic delay was found. This analysis supports locations II and III as potential region of acceleration and excludes location I in Fig. 2.13.

The situation of two type III bursts (labeled A and B in Fig. 2.13) associated with a single spike source, as was found in two of the 92/08/18 events, suggests location II as a possible acceleration region. Position III is only consistent if the field lines meet in position II and continue in parallel to position III.

### 2.6 Conclusions

In all analyzed events the spike sources are always located at positions coinciding with expected locations from extrapolated type III trajectories to lower altitudes. These observations thus strongly support a model for radio spikes occurring in the course of type III beam propagation or near its origin, consistent with independent spectrogram observations (Benz et al. 1996). They add further evidence for spikes being a signature of the mechanism accelerating electron beams that cause type III bursts. This appears to be the simplest interpretation (cf. Fig. 2.13).

This property is supported by the results of the 92/08/18 observations, where in two events simultaneous or consecutive type III bursts on different magnetic field lines originated in the same spike source. Energetic electrons appear to be injected into different and diverging coronal structures from one single position. Such a diverging magnetic field geometry is the standard ingredient of reconnection. These observations are consistent with the hypothesis that metric spikes may be a signature of particle acceleration.

Earlier imaging investigations on metric spikes associated with type III bursts and their interpretation (Krucker et al., 1995, 1997) can now be compared to these additional observations. They proposed a scenario of energy release at high altitude with up- and downward moving energized electrons. The upward moving electrons produce type III bursts while propagating along open field lines and the downward moving part loses its energy to the lower corona, transition region or upper chromosphere. The radio emission of electron beams moving downward from coronal acceleration sites has occasionally been detected (e.g. Klein et al., 1997), but no signature was seen in the observations presented here. High-sensitivity observations are necessary to investigate the processes below the spike source, i.e. above the spike frequency band.
Chapter 3

Electron Firehose Instability and Acceleration of Electrons in Solar Flares

X-rays will prove to be a hoax.

Lord Kelvin (1824–1907)

**ABSTRACT:** An electron distribution with a temperature anisotropy \( T_\perp < T_\parallel \) can lead to the Electron Firehose instability (Here \( \parallel \) and \( \perp \) denote directions relative to the background magnetic field \( B_0 \)). Most of the acceleration mechanisms that can account for the energization of the bulk of electrons to X-ray producing energies of \( \sim 20 \) keV exhibit a preference of accelerating particles in parallel direction. Electron velocities therefore become more and more aligned along the background magnetic field in the course of the acceleration process and an anisotropy in the electron velocity distribution with \( T_\perp < T_\parallel \) is expected during the impulsive phase of a flare. The properties of the excited waves and the thresholds for instability are investigated by using linearized kinetic theory. These thresholds were connected to the pre-flare plasma parameters by assuming an acceleration model acting exclusively in parallel direction. For usually assumed pre-flare plasma conditions the electrons become unstable during the acceleration process and left-hand circularly polarized waves with frequencies of about \( \sim |\Omega_p| \) are excited at parallel propagation. Indications have been found, that the largest growth rates occur at oblique propagation and the according frequencies lie well above the proton gyrofrequency.

3.1 Introduction

Particle acceleration is a phenomenon occurring at many different sites throughout the universe. An important example of particle acceleration are solar flares, offering a wide range of observations that allow one to probe electron and ion acceleration. It is now
widely accepted that the hard X-ray emission observed by various spacecraft reflects
the energization of almost all electrons in the flaring plasma to energies up to \( \sim 25\text{keV} \).
These observations and the observed magnetic fields encompassing the solar flare suggest
that most of the dissipated energy is released by restructuring the magnetic field, e.g.
magnetic reconnection events.

During the impulsive phase of the flare, when the most powerful energization takes
place, electrons must be accelerated to mean energies of about \( \sim 25\text{keV} \) at a rate of
about \( 10^{36} \) electrons per second in order to sustain the observed intensity of the hard
X-ray bursts. Taking the impulsive phase of a flare to last about 10s and assuming an
electron density of about \( 10^{10}\text{cm}^{-3} \) (Moore & Fung, 1972; Vaiana & Rosner, 1978), the
bulk energization must process a coronal volume of at least \( 10^{27}\text{cm}^3 \). Thus energization
must affect a large fraction of the electron population in the flaring region.

In view of this background we want to briefly describe the processes which may be
responsible for particle acceleration in solar flares. For a detailed review of possible
acceleration processes in impulsive solar flares see e.g. Miller et al. (1997) and Cargill
(1999).

1) Shock Acceleration: There are two types of shock acceleration. The one referred to
as 'shock drift acceleration' involves the shock electric field that reflects and accelerates
the particles moving along the shock surface. Since this mechanism is only effective
when the shock normal approaches an angle of 90° to the background magnetic field,
either the gained energy or the particle flux is very limited. It seems to be unlikely
that this mechanism is responsible for the large number of accelerated particles in solar
flares. The second kind of shock acceleration is called 'diffusive shock acceleration'. In
this process the particles cross the shock-front several times, interacting with scattering
centers on both sides of the shock. In the rest frame of the shock these centers approach
each other and the particles systematically gain energy. This kind of acceleration process
requires a certain initial velocity in order to become effective. The ion velocity has to
exceed the Alfvén speed \( V_A = B_0/\sqrt{4\pi \rho} \) while the electrons must have velocities at least
above \( \sqrt{m_i/m_e} V_A \). This has been called the 'injection problem'.

2) Acceleration by parallel electric fields: Direct acceleration by electric fields depend
on its strength compared to the Dreicer field \( E_D = (2\pi e^3 n_e \ln \Lambda)/(k_BT) \). If \( E > E_D \)
most electrons and ions gain energy. If \( E < E_D \) only electrons in the high energy tail of
the velocity distribution function will be accelerated. The limitation in both cases is the
maintenance of overall neutrality of charge and pre-existing current in the acceleration
region.

3) MHD turbulence: This acceleration mechanism occurs when particles interact
many times with randomly moving MHD waves. Due to a slight over-plus in head-on
collisions the interaction results in an energy gain for the particle. As in the shock accel¬
eration model, the acceleration by MHD turbulence suffers from an 'injection problem':
Thermal ions and electrons cannot resonate with MHD waves for typical solar pre-flare
conditions.

A solution for this problem is the assumption of MHD turbulent cascades that channel
the energy residing in the MHD turbulence to smaller scales and into the region
where interaction with thermal particles is possible. A realization of this scenario is
proposed in Miller (1991), Miller & Roberts (1995) and Miller (1997). An MHD turbulent cascade transfers the energy from large scale MHD waves to smaller scales where the energy may be absorbed by the particles. The mechanism that dissipates the wave energy into the particles is transit-time damping (Fisk, 1976; Stix, 1992). It is basically a resonant Fermi acceleration of second order. Only particles in resonance with a low-amplitude MHD wave are affected. The resonance condition is the usual \( I = 0 \) (or Landau) resonance given by \( \omega - k_v v_l \approx 0 \). As the particles can only gain energy in the direction parallel to the background magnetic field, the temperature in parallel direction increases.

A preference for acceleration along the background magnetic field is a common feature of the acceleration models mentioned above. The velocity distribution thus becomes more and more anisotropic during acceleration. If energization in parallel direction is from a thermal level of some 0.1 keV to 20 keV or more but the perpendicular temperature remains constant, the anisotropy is substantial. The free energy residing in parallel direction may give rise to growth of plasma waves.

For \( T^e_{\parallel} > T^e_{\perp} \) and high beta plasmas, Hollweg & Völk (1970) and Pilipp & Völk (1971) have proposed the Electron Firehose instability. This instability is an extension to higher frequencies of the (MHD) Firehose instability, originally mentioned by Parker (1958). While the Firehose instability is of a completely non-resonant nature, the Electron Firehose instability involves non-resonant electrons but resonant protons.

For large anisotropy of the electron distribution, the electrons become also resonant. This instability is described in Pilipp & Benz (1977) and is called the Resonant Electron Firehose instability.

Having been applied to a variety of problems, the Electron Firehose instability has not been considered to occur during electron acceleration in solar flares.

Here we investigate the threshold for growth of plasma modes resulting from acceleration and infer a prediction for the evolution of the distribution function in velocity space with respect to conditions expected to occur in solar flares.

We assume that no significant instability of Langmuir waves occurs. This is suggested by the following argument: If a large fraction of the available energy normally did go into Langmuir waves, we would expect always a radio signature orders of magnitude higher than ever observed during the impulsive phase (Benz & Smith, 1987).

Section 3.2 describes the techniques used to solve the dispersion equations. In the following section the results obtained are shown and the thresholds of the instability are presented. In Sect. 3.4 we discuss the effect on the acceleration of electrons and conclude this work.

### 3.2 Method

Linearized kinetic theory is applied to the problem. The highly non-linear dispersion equation obtained from this theory is solved by using IDLWhamp, a graphical user interface to the well-known WHAMP code, originally programmed by Rönnmark (1982).
A detailed description of the WHAMP code and the newly programmed IDL interface can be found in Appendix A.

It is assumed in the following that the electron velocity distribution function can be described by a bi-maxwellian with different temperatures in parallel and perpendicular direction with respect to the background magnetic field. Hence for our problem the parameters $\Delta_j$ and $\alpha_j$ were set to unity in equation (A.4).

Taking into account the uncertainties in the acceleration region including a possible pre-heating mechanism altering the pre-flare plasma conditions, we do not want to restrict our work to the parameters of a particular scenario. According to Pallavicini et al. (1977) reasonable ranges in the acceleration region of an impulsive solar flare would be $\approx 100 - 500$G for the background magnetic field, $\approx 10^9 - 10^{11}$cm$^{-3}$ for the number density and $\approx 10^6 - 10^7$K for the temperature of electrons and protons.

### 3.3 Results

#### 3.3.1 Electron Firehose Instability

![Dispersion Relation Diagram](image)

**Figure 3.1:** A typical plot of the dispersion relation. The chosen parameters are $T^e_\parallel = T^p_\parallel = T^p_\perp = 10^7$K, $T^e_\parallel/T^e_\perp = 20, n_e = 5 \cdot 10^{10}$cm$^{-3}, B_0 = 100$G. The real part of the frequency $\omega_r$ and the growth rate $\gamma$ are normalized to the proton gyrofrequency $|\Omega_p|$. The parallel wave vector is normalized to the proton inertial length. The whole branch is left-hand circularly polarized.

The only mode exhibiting significant growth rates in our calculations is a left-hand circularly polarized wave which was identified to be the Electron Firehose instability.
This mode evolves out of a stable right-hand polarized whistler wave at small anisotropy. With increasing anisotropy, the frequency \( \omega_r \) is shifted so that, with the convention \( \omega_r > 0 \), the mode becomes left-hand circularly polarized at \( k \parallel B_0 \) in the unstable regime, cf. section 7 in Gary (1993). A typical dispersion relation is plotted in Fig. 3.1.

![Figure 3.1: Dispersion relation for the left-hand polarized whistler wave.](image)

**Figure 3.2:** Resonance factor \( |\zeta^-| \) of the protons versus the electron anisotropy \( T_\parallel/T_\perp \) for the fastest growing modes. The larger \( |\zeta^-| \), the smaller is the fraction of protons in resonance.

By introducing the resonance factor

\[
\zeta_j^\pm = \frac{\omega_r \pm \Omega_j}{\sqrt{2}|k||v_{j\text{th}}|}, \quad (3.1)
\]

the values for the protons and electrons are found to be \( |\zeta_p^-| \sim 1 \) and \( |\zeta_e^-| \gg 1 \), demonstrating resonance for the protons and non-resonance for the electrons.

According to Hollweg & Völk (1970) there are also right-hand circularly polarized modes, which can become unstable for this extension of the Firehose instability. These modes have been found, but the growth rates are smaller than the growth rates of the left-hand polarized modes described above.

As the instability first appears, the phase velocities of the resonant waves are near the peak of the proton distribution. Fig. 3.2 displays a plot of the proton resonance factor (3.1) for the fastest growing modes versus electron anisotropy. With increasing anisotropy less protons are resonant and the resonance factor increases. The change in the fraction of resonant protons is mirrored in the excited frequency range. As depicted in Fig. 3.3 the unstable frequency range grows to a maximum value at an anisotropy of about \( T_\parallel/T_\perp \sim 12 \), coinciding with the maximum value of the resonance factor at \( |\zeta_p^-| \sim 0.57 \) of the protons (cf. Fig. 3.2). As the resonance factor decreases
Figure 3.3: The growth rate $\gamma$ versus the real part of the frequency $\omega_r$, both normalized to the proton gyrofrequency, for different anisotropies $T_{||}^e/T_\perp^e$. The values of the other plasma parameters are the same as in Fig. 3.1.

again, the excited frequency range becomes narrow around $\omega_r/|\Omega_p| \sim 1$. This narrow range is in itself evidence for the resonant character of the instability.

### 3.3.2 Instability Threshold

In this section we present the calculated threshold for linear growth of L-mode waves excited by the Electron Firehose instability.

The initial plasma is assumed to be maxwellian with temperatures $T_{\perp0} = T_{||0} = T_0$ perpendicular and parallel to the background magnetic field for both plasma species, the electrons and the protons. Taking into account an acceleration mechanism for the electrons acting only in parallel direction, the perpendicular temperature remains constant throughout the whole acceleration process, i.e. $T_{\perp e} = T_{\perp0}^e$. In order to investigate the condition of the pre-flare plasma for the Electron Firehose instability to occur during the acceleration process, the initial plasma parameters have to be connected to the actual plasma parameters during the acceleration. With the assumptions above, this can be done by defining an initial parallel plasma beta, $\beta_{||0}^e$, via the perpendicular plasma beta

\[ \beta_{||0}^e \equiv \beta_\perp^e = \frac{8\pi n_e k_B T_{\perp}^e}{B_0^2}, \tag{3.2} \]

and the usual parallel plasma beta by

\[ \beta_{||}^e \equiv \frac{8\pi n_e k_B T_{||}^e}{B_0^2} = \beta_{||0}^e \cdot \frac{T_{||}^e}{T_\perp^e}, \tag{3.3} \]
where the connection between these two quantities is given by the temperature anisotropy $T^e/T^\perp$.

![Graph](image)

**Figure 3.4:** The maximum growth rate $\gamma_{\text{max}}$ normalized to the proton gyrofrequency $|\Omega_p|$ versus the anisotropy $T^e/T^\perp$ of the electrons; $T^p = T^e$, $T^\perp = T^\perp$. The appropriate frequency is always of the order of $|\Omega_p|$.

According to Hollweg & Völk (1970) the instability criterion for the Electron Firehose instability may be approximated by

$$1 - \beta_0^e A_e < 0,$$

where the anisotropy factor is defined by $A_e = 1 - T^e/T^\perp$.

As one can see from inequality (3.4), the instability threshold does not depend directly on the parameters $n_e$, $T^e$, $B_0$, but only on the resulting $\beta_0^e$. This independence is also reproduced with the numerically obtained data. For our purpose, the plasma is therefore fully described by the plasma beta.

In Fig. 3.4 the function $\gamma_{\text{max}}(T^e/T^\perp)$ is plotted for five different values of $\beta_0^e$. The maximum growth rate of the instability steeply raises at the threshold of the instability and flattens for larger anisotropies, where $\gamma_{\text{max}}/|\Omega_p|$ approaches unity.

From these results, the contour of zero growth rate in the $A_e - \beta_0^e$ plane has been derived (cf. Fig. 3.5). The discrepancy between the analytically derived relation (3.4) and the numerically obtained values is due to the approximation used in Hollweg & Völk (1970).

According to inequality (3.4), instability cannot occur for a parallel beta smaller than unity. Due to the deviations of the approximation mentioned above, this limit is shifted to a value of $\sim 1.6$ (cf. Fig. 3.5).
Figure 3.5: Threshold for the Electron Firehose instability in anisotropy factor vs. parallel electron beta. The scaling of both axes is logarithmic. The dashed curve shows the instability limit according to equation (3.4). The dotted line represents a fit to the numerically obtained values which are depicted by diamonds. The areas above the respective lines are the unstable regions.

In order to investigate the necessary properties of the pre-flare plasma for the Electron Firehose instability to occur, it is the initial plasma beta that is of interest. Fig. 3.6 depicts the same plot as Fig. 3.5 but this time the anisotropy factor $A_e$ has been plotted versus the initial plasma beta $\beta_{i0}^e$. The dotted line in both figures represents a fit to the numerically obtained values and is an extrapolation to a broader range of beta values. The negative $\beta_{i0}$ at the $A_e \rightarrow 1$ limit is an artifact of this extrapolation.

The values of the initial plasma beta for the Electron Firehose instability to occur at considerable values of $T_{\parallel}^e/T_{\perp}^e$ are well within the range of usually assumed pre-flare plasma parameters. For example, an initial plasma beta of $\beta_{i0}^e \approx 0.05$ can be realized by assuming pre-flare plasma parameters of $n_e = 5 \times 10^{10} \text{cm}^{-3}$, $T_0^e = 3 \times 10^6 \text{K}$ and $B_0 = 100 \text{G}$. This plasma becomes unstable at an anisotropy of $T_{\parallel}^e/T_{\perp}^e \approx 32$.

According to the acceleration model via transit-time damping, this is a reasonable value for the anisotropy to occur during the acceleration process (Lenters & Miller, 1998).

### 3.3.3 Influence of Anisotropic Protons

If we assume the protons to be heated by the same or a similar mechanism, it is to be expected that they will grow anisotropic in the same way the electrons do. Hence, we
also have investigated the influence of anisotropic protons and briefly discuss the effect of an additional proton anisotropy on the instability.

Consider a plasma with anisotropic electrons and isotropic protons that is already unstable to the Electron Firehose instability. When the protons are anisotropized by increasing the temperature in parallel direction, more and more become resonant to the L-waves, non-resonantly excited by the electrons. As shown by Hollweg & Völk (1970), the protons are damping these waves. Hence, it is to be expected that the resulting growth rate of the L-waves decreases as the proton anisotropy is increased. This expectation has been verified by numerical calculation.

Moreover, the protons are heated by absorption of the excited waves at the expense of the electrons (Pilipp & Völk, 1971). According to Kennel & Petschek (1966) this scatters the protons to higher perpendicular velocities and hence, destroys or inverts the parallel proton anisotropy. It inhibits the growth of the electron anisotropy and may complicate the acceleration to higher energies, but increases the bulk energy of the protons. The energy transfer from the electrons to the protons via the Electron Firehose instability could be responsible for the proton energization, which is a problem in the transit-time damping scenario (Miller, 1998).

If the protons are anisotropic, there is an additional right-hand polarized wave mode. This mode is the extension of the left-hand Electron Firehose mode to negative frequencies. According to Hollweg & Völk (1970) this mode has real frequencies mainly below the proton gyrofrequency and is in resonance with the protons. As the anisotropy of the

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**Figure 3.6:** The same plot as Fig. 3.5 but this time \( A_e \) has been plotted vs. the initial electron beta. Again, as in Fig. 3.5, the dashed curve shows the instability limit according to equation (3.4) and the dotted line represents a fit to the numerically obtained values.
Figure 3.7: Frequency and the according growth rates vs. the propagation angle $\Theta$ with respect to the background magnetic field. The dashed lines indicate the branch of the parallel Electron Firehose instability. The solid lines show the oblique mode that exhibit much faster growth than the Electron Firehose instability. The small plot is an enlargement of the region, where the crossing of the growth rates occurs. The plasma parameters are the same as in Fig. 3.1.

As electrons increases, this right-hand polarized mode becomes less and less significant.

The proton cooling through the instability of the right-hand mode competes with the heating by the left-hand mode and it is not yet clear what the net energization of the protons will be.

3.3.4 Oblique Propagation: Preliminary Results

Sample calculations in oblique directions indicate an unstable branch of solutions that grows faster than the parallel Electron Firehose mode. This mode is stable at parallel propagation and is also left-hand circularly polarized. Fig. 3.7 depicts frequency and maximum growth rate versus the angle $\Theta$ of both modes. The dashed lines represent the branch, that is excited by the Electron Firehose instability at parallel propagation. At a propagation angle of about $10^\circ$ with respect to the background magnetic field, the growth rate of the oblique mode becomes larger than the growth rate of the parallel mode. Hence, the determining mode for instability thresholds is the oblique mode rather than the parallel Electron Firehose mode.

As calculations have shown, not only the growth rate of the oblique mode is larger than in the parallel case, but also the instability threshold may be lower with respect to
anisotropy. Plasmas being stable with respect to the parallel mode exhibited instability to the oblique mode. Therefore, the thresholds for instability derived in the section above can be considered as upper boundaries.

Left-hand circularly polarized oblique modes seem to have never been considered in connection with the Electron Firehose instability. The further investigation of these modes is the subject of ongoing work.

### 3.4 Discussion and Conclusion

Numerical solutions of the dispersion equation for left-hand circularly polarized electromagnetic waves, propagating parallel to the background magnetic field, have shown that the Electron Firehose instability, usually considered as a 'high-beta plasma' instability, must be expected in coronal plasmas which are processed by an acceleration mechanism with a preference in parallel direction. The distribution function of the electrons in velocity space has been represented by a bi-maxwellian with temperatures $T_\perp$ and $T_\parallel$, perpendicular and parallel with respect to the background magnetic field. The protons have been assumed to be isotropic and in thermal equilibrium.

Considering the uncertainty in the pre-flare conditions, we have investigated instabilities in a broad range of plasma temperature $T$, density $n$ and background magnetic field $B_0$. For these plasmas, it was found that the Electron Firehose instability occurs at anisotropies that must be expected for an acceleration mechanism acting predominantly in parallel direction and being capable of producing the observed electron bulk energization.

The unstable parallel modes that have been found are left-hand circularly polarized and non-resonantly excited by the electrons, but partially cyclotron resonant with the protons, which absorb the wave energy. Hence, energy is transferred from the electrons to the protons. Assuming the density and the magnetic field to be constant during the acceleration process, there is a limiting electron temperature in parallel direction that cannot be exceeded without loosing energy to the protons via the Electron Firehose instability. The Electron Firehose instability may thus inhibit the electron acceleration process and limit the reachable energies.

At angles $\Theta \neq 0$ with respect to the background magnetic field, an oblique mode has been found, that exhibits even larger growth rate than the parallel Electron Firehose instability. This mode is also left-hand circularly polarized. The properties of the oblique mode open new aspects on the Electron Firehose instability and its thresholds.

Taking anisotropic protons into account, the left-hand mode may extend to negative frequencies. They correspond to the right-hand circularly polarized mode resonantly excited by the protons. Due to the cooling effect of this instability it is not yet clear to what the net energization of the electrons and the protons will amount. This topic will be the subject of future particle simulations.
Chapter 4

Test Particle Simulation of the Electron Firehose Instability

*I think there is a world market for maybe five computers.*

Thomas J. Watson (1874–1956)

**ABSTRACT:** In the course of the energization of the bulk of electrons to energies of some tens of keV during the impulsive phase of a solar flare, the velocity distribution function of the electrons is predicted to become anisotropic with $T_{||}^e > T_{\perp}^e$ (Here, $||$ and $\perp$ denote directions with respect to the background magnetic field). Such a configuration can become unstable to the so-called Electron Firehose instability (EFI). Left-hand circularly polarized electromagnetic waves propagating along the magnetic field are excited via a non-resonant mechanism: electrons non-resonantly excite the waves while the protons are in resonance and carry the wave. The non-resonant nature of the instability raises the question of the response of the electron population to the growing waves. Test particle simulations are carried out to investigate the pitch-angle development of electrons injected to single waves and wave spectra. To interpret the simulation results, a drift kinetic approach is developed. The findings in the case of single wave simulations show the scattering to larger pitch-angles in excellent agreement with the theory. The situation dramatically changes when assuming a spectrum of waves. Stochasticity is detected at small initial parallel velocities resulting in significant deviations from drift kinetic theory. It enhances the scattering rate of electrons with initial parallel velocity below to the mean thermal perpendicular velocity. Increased scattering is also noticed for electrons having initial parallel velocity within an order of magnitude of the resonance velocity. The resulting pitch-angle scattering is proposed to be an important ingredient in Fermi-type electron acceleration models, particularly transit-time acceleration by compressional MHD waves.
4.1 Introduction

A solar flare is a complex explosive phenomenon in the solar corona. Apparently, magnetic energy is abruptly released, driving bulk mass motion, heating plasma and accelerating electrons and ions to hard X-ray energies and beyond (e.g. Sturrock, 1980). Particles are accelerated at many sites throughout the universe but solar flares offer the widest range of observations and allow to probe both, electron and ion acceleration. During solar flares, large amounts of energies ($\sim 10^{28}$ to $\sim 10^{34}$ erg) are released on time scales varying from fractions of seconds to minutes (Kiplinger et al., 1984). The magnitude and energy spectrum in hard X-rays often indicate that much of the released energy goes into the bulk energization of electrons at tens of keV. In addition, observations show that the number of energized electrons are comparable to the electron content of the whole flaring region. The requirement of accelerating such a large number of electrons ($\sim 10^{37}$) within about 10 s (Brown, 1971; Moore & Fung, 1972) restricts the candidates of accelerators. Currently, possible mechanisms can be roughly divided into three broad classes: stochastic acceleration by MHD waves, shock acceleration and acceleration by direct electric fields (dc) (for a review see e.g. Miller et al., 1997).

A common feature of these accelerators is a preference in accelerating particles in the direction parallel to the background magnetic field. Parallel dc electric fields trivially accelerate only along the magnetic field while stochastic scenarios as transit-time damping act via small amplitude magnetic mirroring, which is only capable of transferring energy in parallel direction if no additional scattering mechanism is provided (Lenters & Miller, 1998). Works of e.g. Wu (1984) and Leroy & Mangeney (1984) describe the parallel directed energization of electrons at the earths bow shock via shock drift acceleration with quasi perpendicular shocks. The low collision rate in the solar pre-flaring plasma cannot reduce the significant anisotropy in the velocity distribution function of the according particle species (e.g. electrons) that builds up during the acceleration process. When assuming acceleration from a thermal level up to orders of $\sim 10$ keV, the anisotropy becomes essential.

Some of the acceleration processes mentioned above require pitch-angle scattering in order to operate efficiently, among which is a scenario proposed by Miller et al. (1996) based on transit-time acceleration. The primary energy input is mass motion in form of an outflow with velocities of order of the Alfvén speed resulting from magnetic reconnection events. Due to its Alfvénic speeds, the outflow can excite turbulence and an MHD cascade is initiated. The cascade transfers the energy from the initial large-amplitude and large-scale MHD waves to small scales and small amplitudes where the energy can be dissipated into thermal electrons, by transit-time damping (Fisk, 1976; Stix, 1992). Transit-time acceleration can be regarded as a low-amplitude realization of the Fermi process with the main difference, that the wave-particle interaction is, instead of non-resonant magnetic mirroring, of rather resonant nature, i.e. $v_\parallel \approx \omega/k_\parallel$.

It has been shown (Lenters & Miller, 1998) that this mechanism is in need of a pitch-angle scattering mechanism in order to operate efficiently. The pitch-angle scattering time scale has to be comparable to the characteristic wave-particle interaction time which is given by $\approx \pi/kv_A$, where $v_A$ is the Alfvén speed. Pitch-angle scatter-
ing by Coulomb collisions is therefore not sufficiently rapid to lead to efficient electron acceleration by transit-time damping (Lenters & Miller, 1998). Whistler waves (right-hand circularly polarized waves below \( \Omega_e \)) at a low energy level have been proposed as possible scatterers (Miller, 1997). These waves would deliver the desired scattering, but it is unclear whether they are present and what the actual source would be. Since the bulk energization of electrons in impulsive solar flares is essential, a successful acceleration scenario must contain inherently all necessary ingredients. In order to obtain a self-contained picture of bulk acceleration during solar flares, the pitch-angle scattering mechanism therefore preferably originates within the acceleration scenario. The anisotropy in the electron velocity distribution described in the paragraph above is a source of free energy and therefore may deliver a possible agent to scatter electrons to higher pitch-angle.

It is known from analyses in linearized kinetic theory (Hollweg & Völk, 1970; Pilipp & Völk, 1971) that an anisotropic plasma \( (T_\parallel > T_\perp) \) can become unstable to the so-called Electron Firehose instability (EFI). The EFI is an extension of the well known (MHD-) Firehose instability, originally mentioned and applied the the solar wind to explain the unexpected isotropy in the particle distributions by Parker (1958). While the classic Firehose instability is of entirely non-resonant nature, the EFI involves resonant protons with the electrons remaining non-resonant. It has been proposed in a former work (Paesold & Benz, 1999) that Electron Firehose (EF) waves may well be excited in course of the acceleration processes during solar flares. In addition to the parallel propagating EF mode a new branch of the EFI was found at oblique angles. This mode was further investigated by Li & Habbal (2000). Messmer (2002) showed by using PIC simulations that the parallel EFI is capable to scatter the electrons on short time scales in pitch-angle. He uses large anisotropies to achieve a rapid development. This results in resonant wave-electron interaction, and hence, very efficient pitch-angle scattering is obtained.

The work presented herein focuses on the non-resonant phase of the process. Since the source of free energy lies in the anisotropic distribution function, the instability is expected to isotropize the electron population and to remove the anisotropy. Do electrons just loose parallel energy or are they also scattered in pitch-angle? Which non-resonant processes scatter electrons in velocity? Are all electrons affected equally? Even though the non-resonant scattering is expected to be less effective than a resonant interaction, an understanding of the non-resonant processes leads to important insights into the plasma behavior near the instability threshold. Before reaching large anisotropies, the instability starts eroding it and the plasma eventually resides in a state close to marginal instability. The non-resonant situation at small electron anisotropies, close to the threshold of the EFI, is probably more realistic for solar applications.

Besides the importance of a possible application to solar flare physics, the EFI poses the fundamental problem of a kinetic description for the electron response to a hydro-magnetic instability. The problem is analyzed by drift kinetic theory and by applying test particle simulations. The response of the electrons to the EF wave field, obtained from linearized kinetic theory, is investigated in view of the displacement in velocity space, mostly in pitch-angle.
The plan of the paper is as follows: the linearized kinetic analysis of the EFI is presented in Sect. 4.2. An analytic description of an electron in an EF wave field is presented in Sect. 4.3 while the test particle simulations are described in Sect. 4.4. The presentation of the results in Sect. 4.5 is followed by a discussion in Sect. 4.6 and the work is concluded by the summary in Sect. 4.7.

4.2 Electron Firehose Instability

Figure 4.1: A typical plot of the dispersion relation of EF waves. The chosen parameters are $T_\perp^e = T_\perp^p = T_\parallel^p = 10^7$ K, $T_\parallel^e/T_\perp^e = 13$, $n_e = 5 \cdot 10^{10}$ cm$^{-3}$, $B_0 = 100$ G. The real part of the frequency $\omega_r$ and the growth rate $\gamma$ are normalized to the proton gyrofrequency $\Omega_p > 0$. The parallel wave vector is normalized to the proton inertial length. The whole branch is fully left-hand circularly polarized.

The extension of the (MHD-) Firehose instability to higher frequencies, the Electron Firehose Instability (EFI), is non-resonantly excited by the electrons, whereas the protons are now in resonance with the waves. The anisotropy in the electron velocity distribution drives the waves while the protons carry the wave. A typical dispersion relation of the EF waves is displayed in Fig. 4.1. The mode depicted is parallel propagating, purely transverse and left-hand circularly polarized. The dispersion of the EF waves is computed by IDLWhamp (Paesold, 2002), an easy to use IDL interface to the WHAMP code originally developed by Rönnmark (1982). The code provides the user with the full solution of the dispersion equation in linearized kinetic theory. For simplicity, the anisotropy is modeled by maxwellian distributions having $T_\parallel^e > T_\perp^e$, and only the parallel mode is considered.
While the protons have been assumed to remain isotropic in Fig. 4.1, a proton anisotropy alters the linearized dispersion of the waves. The $k_\parallel$ of maximum growth is shifted towards smaller $k$ with increasing proton anisotropy and and the according frequency $\omega_r$ decreases. The mode can also become right-hand circularly polarized at small values of $k_\parallel$ when the proton anisotropy exceeds a threshold value of $T^p_\parallel / T^p_\perp > 2$. In the analysis presented herein, a possible proton anisotropy is not taken into account.

### 4.3 Analytical Theory

The electron displacement in velocity space is determined by the non-resonant interaction with the wave's electromagnetic field, which is decomposed into its parallel component $E_\parallel = \hat{e}_\parallel \cos \Psi$ and two circularly polarized components $E_\pm = \hat{e}_\pm (\cos \Psi, \mp \sin \Psi)$ with amplitudes $\hat{e}_\pm = (E_x \pm E_y)/2$, where $\Psi = k_\parallel z + k_\perp x - \omega t$ is the wave's phase at the location of the electron. The negative subscript refers to left-hand polarized waves. The magnetic field of the wave is obtained via Faraday's law $\mathbf{B} = -c \nabla \times \mathbf{E}$ as $(c/\omega) \left[ \pm (k_\parallel \hat{e}_\pm) \sin \Psi, (k_\parallel \hat{e}_\pm - k_\perp \hat{e}_\parallel) \cos \Psi, \pm (k_\perp \hat{e}_\pm) \sin \Psi \right]$. The relative phase angle between the electron and the wave is denoted with $\phi_\pm = \theta \pm \Psi$. In the following we assume parallel propagating EF waves and therefore $k_\perp = 0$, $k_\parallel = k$, $\hat{e}_\parallel = 0$ and $\hat{e}_\perp = \hat{e}$.

In order to identify the main contributions responsible for the response of the electrons to the waves, adiabaticity is assumed in the analytical approach. Since the wave frequency is much smaller than the cyclotron frequency of the electron, i.e. $\omega_r \approx \Omega_p \ll \Omega_e$, this assumption holds for parallel particle velocities small enough to ensure large frequency separation of the Doppler shifted wave frequency and the electron gyrofrequency. The gyrofrequencies are defined as $\Omega_\alpha = |q_\alpha| B_0 / cm_\alpha$, where $\alpha$ denotes the particle species. In addition, small field strengths have to be assumed for adiabaticity. With these assumptions a drift kinetic approach can be chosen (Hasegawa, 1975). The according continuity equation in phase space is given by

$$\partial_t f_D + \nabla \cdot (\vec{v}_D f_D) + \frac{\partial}{\partial v_\parallel} \left( \frac{F_\parallel}{m} f_D \right) = 0 , \quad (4.1)$$

where $f_D(v_\parallel, \mu, \vec{x}, t)$ is the distribution function of the guiding centers and $\mu = mv_\parallel^2/2B_0$ is the magnetic moment. $F_\parallel$ refers to a possible external force in parallel direction to the background magnetic field. This equation is obtained by transforming the Vlasov equations to guiding center coordinates and by time averaging over the fast gyration of the electrons. In the following, when speaking of quasi-particles, we refer to the gyro centers of the electrons. The drift velocity $\vec{v}_D$ can be split into parallel and perpendicular components and contains several different contributions as e.g. $E \times B$ drift, curvature drift and polarization drift. Whereas the latter drift is negligible for the situation assumed herein, the relevant drifts are the $E \times B$ and the curvature drift given by

$$\vec{v}_E = \frac{(c\vec{E} + \vec{v}_\parallel \times \vec{B}) \times \vec{B}}{B^2} , \quad (4.2)$$

$$\vec{v}_C = \frac{m_e c v_\parallel^2}{q B^2} \left[ \vec{B} \times \left( \vec{B} \cdot \nabla \right) \vec{B} \right] , \quad (4.3)$$
where the parallel direction refers to $\vec{B}_0$. The modification involving the parallel velocity in the $E \times B$ drift has to be taken into account when assuming additional perpendicular magnetic fields (Hasegawa, 1975). Assuming $\vec{E} = \delta \vec{E}$ and $\vec{B} = \vec{B}_0 + \delta \vec{B}$, where $\delta \vec{E}$ and $\delta \vec{B}$ are the wave components, Eqs. (4.2) and (4.3) yield, keeping only contributions of first order in the perturbation, a total drift in perpendicular direction

$$\vec{v}_{D\perp} = \frac{c\bar{e}B_0}{B^2} \left[ 1 - \frac{v_\parallel}{c} \left( \frac{ck}{\omega} \right) - \frac{m_e v_\parallel^2 \omega B_0}{c e B^2} \left( \frac{ck}{\omega} \right)^2 \right] \left( -\sin \Psi \right) \left( \cos \Psi \right)$$

and accordingly, when taking second order contributions into account, the resulting drift in parallel direction

$$v_{D\parallel} = \frac{c\bar{e}^2}{B^2} \left( \frac{ck}{\omega} \right) \left[ 1 - \frac{v_\parallel}{c} \left( \frac{ck}{\omega} \right) - \frac{m_e v_\parallel^2 \omega B_0}{c e B^2} \left( \frac{ck}{\omega} \right)^2 \right].$$

Equating the expression in the brackets in Eqs. (4.4) and (4.5) to zero, directly yields a critical velocity $v_{\text{crit}}$ in parallel and perpendicular direction where the contributions are in balance and result in zero drift. Expanding the numerator of $v_{\text{crit}}$ for $(B_0/B)^2 \ll \Omega_e/\omega$ yields as the only physical solution

$$v_{\text{crit}} \approx \frac{\omega}{k},$$

the phase velocity of the wave. The drift in parallel direction below the critical velocity $v_{\text{crit}}$ is in positive direction and changes sign above it. The magnitude of the perpendicular drift $v_{D\perp} = |\vec{v}_{D\perp}|$ has a minimum (here equal to zero) at $v_{\text{crit}}$. The explanation is straightforward: since the particle propagates at the same parallel speed as the wave the resulting $E \times B$ drift contributions cancel. There is no relative velocity between wave and particle, hence the particle experiences no curvature drift as well. Note, that the critical velocity does not depend on the wave field strengths.

The resulting quasi-particle pitch-angle can readily be obtained by computing

$$\alpha_c = \arctan \left( \frac{v_{D\perp}}{v_\parallel + v_{D\parallel}} \right).$$

When assuming a wave field of $N$ waves indexed by $i$ with constant amplitudes $\bar{e}_i = \bar{e}$, Eqs. (4.4) and (4.5) have to be modified. In analogy to a random walk in the $v_x - v_y$-plane, the resulting rms drift in perpendicular direction is then given by

$$v_{D\perp} = \sqrt{\sum_{i=1}^{N} \bar{v}_{D\perp,i}^2},$$

where $\bar{v}_{D\perp,i}$ is the drift contribution given by Eq. (4.4) for each wave $i$. It has been assumed, that $N$ is a large number and that the time the particle stays in the wave field
is long enough to sample the largest wavelength in the spectrum. The average parallel drift is obtained as

\[ v_{D||} = \frac{c e^2}{B^2} \sum_{i=1}^{N} \left( \frac{c k_i}{\omega_i} \right) \left[ 1 - \frac{v_{\parallel}}{c} \left( \frac{c k_i}{\omega_i} \right) - \frac{m c v_{\parallel}^2 \omega_i B_0}{c e B^2} \left( \frac{c k_i}{\omega_i} \right)^2 \right] . \quad (4.9) \]

The squared magnitude of the total magnetic field \( B^2 = |\vec{B}|^2 \) is given by \( B^2 = B_0^2 + \sum (c e k_i / \omega_i)^2 \). The drift velocities in Eq. (4.7) have to be adapted accordingly to obtain the pitch-angles of the particles in the wave field.

In the case of a wave spectrum, a critical velocity cannot be defined as easily as in the single wave case. An estimate for \( v_{\text{crit}} \) in case of a wave field can be obtained by determining the minimum in Eq. (4.8) or the zero of Eq. (4.9) in dependence on \( v_{\parallel} \).

An important quantity has been omitted in the derivation above. Contrary to the single wave case, the modulus of the total instantaneous magnetic field of the wave spectrum changes as the particle propagates. Since the motion is assumed to be adiabatic, the particle’s magnetic momentum is an invariant. When changing the modulus of the magnetic field by \( \Delta |B| \), the resulting shift in parallel velocity is approximately given by

\[ \Delta v_{\parallel} \approx -\frac{1}{2} \frac{v_{\parallel}^0}{v_{\parallel}^0} \frac{\Delta |B|}{|B_0|} v_{\perp}^0 . \quad (4.10) \]

Here it is assumed that \( \Delta v_{\perp} \ll v_{\perp}^0, \Delta |B| \ll |B_0| \) and that energy is approximately conserved on the average. The first requirement is only true for electrons with non-vanishing initial perpendicular speeds. It is not possible to include \( v_{\parallel}^0 \) to the approximation since \( \Delta |B| \) varies individually for each particle and depends on the relative phases of the applied wave fields. Invariance of the magnetic moment causes a random scatter of the particles and can qualitatively explain the deviation of the simulation results presented in Sect. 4.5.2 from the predictions of Eqs. (4.7), (4.8) and (4.9).

### 4.4 Test Particle Simulation

#### 4.4.1 Simulation Setup

The relativistic equations of motion for a particle of charge \( q \) and mass \( m \) in a field of \( N \) waves and a homogeneous background magnetic field \( B_0 \) are given by

\[
\frac{d\vec{\sigma}}{dt} = q \frac{\vec{\sigma}}{c} \times \vec{B}_0 + q \sum_{i=1}^{N} \left( \vec{E}_i + \frac{\vec{\sigma}}{c} \times \vec{B}_i \right) , \quad (4.11)
\]

\[
\frac{d\vec{x}}{dt} = \vec{\sigma} , \quad (4.12)
\]

where \( \vec{x} \) is the particle position vector, \( \vec{\sigma} \) the velocity and \( \vec{E}_i \) and \( \vec{B}_i \) are the electric and magnetic fields of the wave \( i \) propagating parallel to \( B_0 \).
Test particle trajectories are calculated by integrating a dimensionless form of Eqs. (4.11) and (4.12) with a standard leap-frog mover following Birdsall & Langdon (1991) (at half time step: acceleration with an extrapolated electric field, rotation around the instantaneous magnetic field, including the wave field and acceleration with updated velocities). The code used herein is a leap-frog version of the code used by Miller & Viñas (1993) and has been tested against their version. The results of both codes are in excellent agreement on time scales relevant for the analysis herein.

The background magnetic field is considered to be homogeneous. The wave frequencies and wave numbers are obtained from the IDLWhamp code (Paesold, 2002). The electric field strength $E_i$ of one single wave $i$ is treated as a free parameter, and the magnetic field components are calculated according to Faraday's law $\vec{B} = -c\vec{\nabla} \times \vec{E}$. The simulations are split into two parts. First, the trajectories of single particles in a single wave have been computed. The results of these simulations are presented in Sect. 4.5.1. During a second set of simulations, groups of 500 particles have been followed in a spectrum of 500 waves. The results of this second set of simulations are presented in Sect. 4.5.2. The 500 monochromatic purely transverse waves are confined to a range of frequencies where the linear growth rate does not drop below $\sim 70\%$ of the maximum growth rate. The applied wave spectrum corresponds to the spectrum depicted in Fig. 4.1 restricted to the range of $\sim 4.1 - 6.5$ kc/ωp.

It is assumed, that the electric field amplitude, normalized to the background magnetic field, of each wave is the same and its numerical value is treated as a free parameter. The wave fields are not present at the beginning of the simulation period. The fields are switched on adiabatically, linearly increasing over a time of $t = 1/\Omega_p$ until they reach the desired value. In the course of the analysis it turned out that, as soon as non-vanishing initial perpendicular velocities are taken into account, it is crucial for the final result to let the waves grow slow enough to ensure adiabaticity of the process. The plasma parameters throughout all simulations are $n_e = n_p = 5 \cdot 10^{10}$ cm$^{-3}$, $T_e^\parallel = T_e^\perp = T_p^\parallel = T_p^\perp = 1 \cdot 10^7$ K, $T_e^\parallel / T_e^\perp = 13$ (see Fig. 4.1) and the background magnetic field is $B_0 = 100$ G. The anisotropy of the plasma is chosen such that the resulting wave field represents the situation close to instability threshold.

Since the wave frequencies are at about the gyrofrequency of the protons (see Fig. 4.1), the simulation time span has to cover at least a few tens of $\Omega_p^{-1}$ in order to sample several gyro periods of the waves. The time scale of interest for the electron motion lies at the inverse gyrofrequency of the electrons. The time steps of the simulation therefore have to be chosen to resolve the gyro period of the electrons. Since the real mass ratio $m_p/m_e \approx 1836$ is used, about $10^4$ iterations per proton gyration have to be computed, resulting in a large computing effort.

A group of electrons, which initially can be described by a density peak in velocity space

$$f(v_\parallel, v_\perp, t_0) \propto \delta(v_\parallel - v_\parallel^0) \times \delta(v_\perp - v_\perp^0),$$

(4.13)
is injected at $(v_\parallel^0, v_\perp^0)$ and is followed in time. Randomization of the electron gyro phases is achieved by choosing the initial parallel coordinate to randomly vary along $\vec{B}_0$. It is
Figure 4.2: Power spectrum of x-component of the normalized momentum. The dashed line indicates the expected gyrofrequency of the electron. The group of peaks to the left represents the wave field, shifted to higher frequencies due to the parallel motion of the particle.

assumed that all of the particles of one group experience the same wave field, i.e. the relative phases of the waves are not changed for one set of \((v_\parallel^0, v_\perp^0)\). For each set of initial velocities, an ensemble of 500 electrons is simulated during a time of \(t \times \Omega_p = 40\).

4.4.2 Simulation Analysis

In order to analyze the motion of the gyro center in the test particle simulation the gyro motion of the electrons has to be eliminated from the full particle motion. This is done by a Fast Fourier transformation (FFT) separating the different contributions in the time series of the components of space coordinates and momenta. The spectrum of a representative run is displayed in Fig. 4.2.

The peak at high frequencies is due to the gyro motion of the electron as indicated by coincidence with the electron gyrofrequency \(\Omega_e\). The peaks at lower frequencies represent the spectrum of the applied wave field. It is Doppler shifted to higher frequencies \(kv_\parallel + \omega_k\) since the spectrum is taken in the parallel co-moving guiding center frame. At high enough parallel velocities, the spectrum of the wave field in Fig. 4.2 would be shifted to the gyroation peak resulting in wave-particle resonance. This situation is excluded from this work by assuming initial parallel velocities small enough to ensure non-resonance.

A portion of an exemplary particle trajectory in the x-y plane is displayed in Fig. 4.3. The gyration of the electron forces it into a fast spiral motion along the local magnetic field. In addition, its gyro center moves in a second, more irregular spiral. This second
motion is much slower and superimposed on the unperturbed electron orbit. The gyro centers can be treated as unmagnetized quasi-particles. If not mentioned otherwise, we always refer to these quasi-particles when describing particle properties in the further analysis of the simulations.

## 4.5 Results

### 4.5.1 Single Wave Simulation

Simulations in the field of a single wave are presented in order to establish the range of validity of the approximations derived in Sect. 4.3. Single electrons at several $v_0^\parallel$ are injected at random relative phases $\phi = \theta - \Psi$ into the field of a single wave with the properties $\omega / \Omega_p = 1$ and $kc / \omega_p = 4.33$, according to the linear dispersion of EF waves (see Fig. 4.1).

At vanishing initial perpendicular velocities $v_0^\perp$, the predictions of the theory and the results of the simulation are in excellent agreement (Figure 4.4). The slight deviation towards large $v_0^\parallel$ is a result of loss of adiabaticity as the Doppler shifted frequency of the wave approaches the gyrofrequency of the electron in its parallel co-moving system. At particle speeds of $v_0^\parallel = \omega / k$ a minimum in pitch-angle is found. The minimum is exactly located at $v_{\text{crit}}$ as predicted by Eq. (4.6). The critical velocity $v_{\text{crit}}$ is indicated by a vertical dashed line in the plot.

The dependence of the pitch-angle increase $\alpha_c$ on the wave field strength $\hat{e}$ is shown
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Figure 4.4: Pitch-angle increase of quasi-particles at several initial parallel velocities \( v_0^0 \) by the interaction with a single wave. The initial perpendicular velocity \( v_\perp^0 \) equals zero. The solid line represents the results obtained from Eqs. (4.4) and (4.5) whereas the squares confer to the results of the simulation. The electric field strength of the wave is \( \hat{\epsilon}/B_0 = 10^{-5} \).

in Fig. 4.5. The theoretical predictions are compared with results of the simulations at three different initial parallel velocities \( v_0^0 \). Simulation results for vanishing initial perpendicular velocity fall all very close to the values predicted by the analytical drift approximation. The theoretically approximated and simulated values show excellent agreement over 4 orders of magnitude up to magnetic wave field strengths of order \( 10 \times B_0 \). Although these values are unreasonably high when dealing with waves resulting from linearized kinetic theory, the results clearly show, that the drift approximation is valid on all scales of field strength as long as the initial gyro motion of the electron vanishes.

The situation is different if \( v_\perp^0 \neq 0 \) is assumed. Equation (4.7) would predict that the resulting \( \alpha_c \) is independent of \( v_\perp^0 \). As can be seen in Fig. 4.5, this is only true for values of the electric wave field below \( \hat{\epsilon}/B_0 \approx 10^{-4} \). Larger fields cause loss of adiabaticity and the approximation therefore becomes invalid. The situation differs depending on the value of \( v_\perp^0 \): adiabaticity is lost first at small values (Fig. 4.5, triangles) when increasing \( \hat{\epsilon} \). At larger \( v_\perp^0 \) values the analytic description yields a good approximation even for large values of \( \hat{\epsilon} \) (Fig. 4.5, crosses).

The upper boundary for the analytical approximation to be valid is at an electric wave field of \( \hat{\epsilon}/B_0 \approx 10^{-4} \). This corresponds to a wave magnetic field of 1.3% of the background magnetic field in the case of EF waves. When increasing the field strength beyond this value, the resulting drifts and displacements in velocity space cannot be
Figure 4.5: Plot of pitch-angle increase vs. electric wave field. The symbols here refer to different initial parallel velocities indicated in Fig. 4.4 by vertical dotted lines: triangle $V_0^0 = 10^7 \text{ cm/s}$, squares $V_0^0 = 10^8 \text{ cm/s}$ and cross $V_0^0 = 10^9 \text{ cm/s}$. The dashed curve shows the results from Eq. (4.7). The symbols close to the dashed curves are the results of the simulations for $V_0^0 = 0$. Solid lines connect the results obtained from the simulations for $V_0^0 = 1.5 \cdot v_{th}$.

predicted anymore by Eqs. (4.4) and (4.5). According to Fig. 4.5 this corresponds to a maximum pitch-angle displacement of about 10 deg that can be described by the approximation.

4.5.2 Wave Spectrum Simulation

When assuming a spectrum of EF waves, several aspects of the model change. Figure 4.6 is a part of the trajectory of a quasi-particle that started at $V_0^0 = 0$ and $V_0^0$ as indicated by the dashed line. As can be seen in Fig. 4.6 the trajectory of the quasi-particles in velocity space becomes irregular and takes excursions to negative parallel velocities. The most important reason is the influence of the first adiabatic invariant, the conservation of the magnetic moment $\mu$, described in Sect. 4.3, Eq. (4.10). Contrary to the case of the single wave simulation, the magnitude of the total magnetic field is a function of time and space and therefore redistributes the perpendicular and parallel velocities. The most important aspect is the influence of the first adiabatic invariant, the conservation of the magnetic moment $\mu$, described in Sect. 4.3, Eq. (4.10). Contrary to the case of the single wave simulation, the magnitude of the total magnetic field is a function of time and space and therefore redistributes the perpendicular and parallel velocities. Combined with the long-term conservation of energy, the invariance of $\mu$ causes a stochastic element in the simulation. The resulting displacement in velocity space depends on the time history of $\vec{B}$, i.e. the total change in $|\vec{B}|$ a particle experiences. This only depends on the initial conditions, i.e. the relative phases in the wave field and the relative phase of the electron, and is therefore random for each particle. Another aspect of $|\vec{B}|$ being a function of time is that the range of validity of the approximation is not as clearly defined as in the case
of a single wave. Even though the $rms$-value of the total perpendicular electric wave field lies well in the valid range, the instantaneous electric field can take an excursion far above the limit. Particles experiencing such excursions do not behave adiabatically anymore and severely alter the statistics of the ensemble.

In the case of very weak wave field strengths, the drift approximation still yields a good prediction of the pitch-angle behavior of an ensemble of electrons in a spectrum of EF waves. The results of such simulations is shown in Fig. 4.7a. The electric field strength of a single wave is $\hat{\epsilon}_k/B_0 = 10^{-6}$, corresponding to a $rms$ value of the total perpendicular electric field of $E_{rms}/B_0 \approx 2.2 \cdot 10^{-5}$. The symbols here and throughout the rest of this section refer to different perpendicular initial velocities as follows: diamond $v_\perp^0 = 0$, triangle $v_\perp^0 = 0.5 \cdot v_{th}$, cross $v_\perp^0 = v_{th}$ and asterisk $v_\perp^0 = 1.5 \cdot v_{th}$. Since the results of the simulation do not depend on $v_\perp^0$ at higher $v_\parallel^0$, only a few representative values are plotted. At higher parallel velocities Eqs. (4.7), (4.8) and (4.9) well approximate the simulated data. The simulation deviates from the theoretically approximated values at smaller $v_\parallel^0$ and the resulting pitch-angles exhibit a dependence on the initial perpendicular speed.

The deviation becomes more and more significant when increasing the electric wave fields. A simulation with strong wave fields of $\hat{\epsilon}_k/B_0 = 10^{-5}$, corresponding to a total perpendicular electric field of $E_{rms}/B_0 \approx 2.2 \cdot 10^{-4}$, is depicted in Fig. 4.7b. The simulation exhibits a behavior completely different from what is expected by the drift approximation. Only at the largest initial parallel velocities the simulation approaches the analytical approximation. The mean pitch-angle shift by interaction with the wave

**Figure 4.6:** Trajectory of a quasi particle in a EF wave spectrum with a $rms$-field of $E_{rms}/B_0 \sim 2.2 \cdot 10^{-5}$. The vertical dashed line indicates the parallel initial velocity.
Figure 4.7: Mean pitch-angle $\alpha_c$ due to interaction with a spectrum of EF waves. Sets of 500 electrons at several initial velocities $(v_0^0, v_0^b)$ have been simulated. The initial values for the gyro centers are thus given by $(v_0^0, 0)$. The dashed curve represents the analytic results obtained from Eqs. (4.7), (4.8) and (4.9) whereas the symbols confer to the results of the simulation for different $v_0^0$ (notation see text). Top: The rms electric field amplitude of the total perpendicular wave field is $E_{rms}/B_0 \approx 2.2 \cdot 10^{-5}$. Bottom: $E_{rms}/B_0 \approx 2.2 \cdot 10^{-4}$. 
field does not significantly depend on the initial perpendicular speed $v_\perp^0$. The value of $v_\parallel$ below which the simulation breaks away from the qualitative behavior expected from the drift approximation linearly depends on the wave field strength. Additional simulations in the range up to $\hat{e}_k/B_0 = 10^{-5}$ not shown here have been carried out to confirm this dependence.

![Figure 4.8: Change in the mean pitch-angle α of the real electrons by interaction with a spectrum of EF waves vs. the initial parallel velocity $v_\parallel^0$. The symbols are the same as in Fig. 4.7. Solid lines represent the average values $\langle \Delta\alpha \rangle$, whereas the dashed lines indicate the distribution $\langle \Delta\alpha \rangle + \sigma$ for the simulations displayed in Fig 4.7b. For comparison, vertical dotted lines indicate the perpendicular and parallel thermal velocity of the electron distribution.](image)

Having investigated the motion of the gyro centers, the situation of the real electrons has to be studied now. The full trajectories of the same simulations as above are analyzed and the resulting displacement of the real electrons in pitch-angle is shown in Fig. 4.8. The change in pitch-angle is shown for simulations with an $rms$-value of the total perpendicular electric field of $E_{rms}/B_0 \approx 2.2 \cdot 10^{-4}$, corresponding to Fig. 4.7b. The dashed curves refer to the values of one standard deviation added to the mean. These values are positive throughout the whole plot, as the distribution is broad. Note: the pitch-angle distribution is not normally distributed.

### 4.6 Discussion

The discussion of the results presented in the preceding section are split into two parts. First, the single wave simulation results are discussed separately and in short while the case of the wave spectrum needs to be addressed more extensively.
In the case of the single wave simulation, the findings are in good agreement in a limited range with the expectations derived from the analytic drift kinetic approach discussed in Sect. 4.3. Assuming vanishing perpendicular initial velocities, the predictions of the theory are in excellent agreement with the simulations on all scales of wave field strengths. The assumption of finite values of $v_\perp^0$ yields a limit in the range of validity of the drift approximation: At low wave fields no difference between finite and vanishing $v_\perp^0$ is observed. By increasing the wave amplitudes above a value of $e_k/B_0 \approx 10^{-4}$, the approximation breaks down (Figure 4.5). The deviation becomes significant first at the lowest initial parallel speeds. The value $e_k/B_0 \approx 10^{-4}$ corresponds to a wave magnetic field of $\approx 13\%$ of the background magnetic field and is concluded to be the limiting field strength for the assumption of adiabaticity in the investigated system.

The case of the wave spectrum simulation requires more extensive discussion than the above. Many effects are observed resulting in deviations from the predictions of the drift approximation. As can be seen in Fig. 4.7, the drift approximation deviates from the simulation first at small initial parallel velocities. When increasing the wave fields, the point where the approximation breaks away from the simulation is shifted to larger values of $v_\parallel^0$. These observations are discussed in detail in the following and possible causes of the deviations are presented.

The major aspect, that has changed in the transition from a single wave to a spectrum of waves, is the time dependence of $|B|$. While in the case of a single wave the total magnetic field of the wave is constant in time, the superposition of the single waves in the spectrum causes the total magnetic field to vary. Several new effects result from this. As mentioned in Sect. 4.3, the change in $|B|$ with time results in additional shifts in velocity space due to the conservation of the magnetic moment (Equation 4.10). The particles are shifted to smaller parallel velocities when the value $|B|$ increases and, if the change is large enough, can therefore be displaced to even negative parallel speeds. This causes large excursions in pitch-angle and severely alters the statistics of the ensemble. The mean is shifted to larger values than expected from the approximation. The conservation of magnetic moment surely influences the simulation but is not included in the analytical approximation. Yet, it cannot account for all effects in the simulation. The significant dependence on $v_\perp^0$ that is predicted by Eq. (4.10) is not observed in the simulation (see Fig. 4.7b). There is no systematic dependence on the initial perpendicular speed in the simulation.

The above assumes that the change in $|B|$ is such that adiabaticity is conserved and, hence, the magnetic moment is conserved. Taking the electric wave field as the responsible parameter, it is known from the single wave simulations that the approximation breaks down at values of about $e_k/B_0 \approx 10^{-4}$. While the rms-value of the total electric field also lies around that value for the simulation presented in Fig. 4.7b, the momentary field can take rather large excursions and typically reach two to three times the rms-value. This causes loss of adiabaticity and therefore results in pitch-angle scattering that cannot be explained by the drift approximation. As can be seen from the single wave simulation in Fig. 4.5, adiabaticity is lost first at small initial parallel speeds while the approximation holds for large $v_\parallel^0$. The largest deviations from the theoretically approximated values occur at small initial parallel velocities. At larger speeds the
approximation is still valid within a 15% range. Due to the faster motion the particle feels the excursions less than the slow particles.

Figure 4.9: Poincaré map of an ensemble of 70 quasi-particles with initial momenta indicated by the vertical bar near (0,0). The initial parallel speed is \( v_0^\parallel = 1 \times 10^7 \) cm/s. The surface in phase space is defined by equating the phase of the particles to zero, i.e. \( \theta = 0 \). The vertical dashed line indicates the position of the minimum in the gyro center pitch-angle \( \alpha_c \) at \( v_{\text{crit}} \) obtained from Eqs. (4.7), (4.8) and (4.9).

To illustrate the stochastic behavior in the simulation, an exemplary Poincaré map for a small ensemble of 70 quasi-particles is presented in Fig. 4.9. The procedure follows Karimabadi et al. (1989). The transition to canonical coordinates yields \( \vec{F}_\parallel = \vec{p}_\parallel \) and \( \vec{F}_\perp = \vec{p}_\perp + q/c \cdot \vec{A} \) for parallel and perpendicular momentum, respectively; \( \vec{p}_\parallel \) refers to the physical momentum and \( \vec{A} \) is the vector potential defined as \( (c/\omega)(\vec{e} \sin \Psi, -\vec{e} \cos \Psi + xB_0, 0) \) according to the definition of \( \delta \vec{B} \) and \( \vec{B}_0 \) in Sect. 4.3. A surface in the phase space is defined by equating the phase angle \( \theta \) of the gyro centers to zero. The dots in the plot therefore represent the phase angle \( \theta = 0 \) for all particles throughout the temporal history of the simulation. The thick black bar indicates the initial values of the real electrons. The finite width in perpendicular direction is a result of the \( xB_0 \) term in the y-component of the vector potential \( \vec{A} \).

The particles fill an area in phase space that is much larger than expected from the drift motion derived in Sect. 4.3. In particular it can be seen that a significant fraction of particles is scattered to negative parallel velocities. A new feature becomes apparent which has not been addressed yet: the particles concentrate and become over-dense in the region around \( v_{\text{crit}} \) (Figure 4.9). Instead of systematically interacting with the wave and experiencing the helical motion in the wave's magnetic field, the particles
feel a randomly fluctuating field at $v_{\text{crit}}$. This results in anomalous stochastic increase in perpendicular velocity. By investigating Poincaré maps at several parallel initial velocities, it turns out that the point where the approximation breaks away from the simulation is determined by the size of the attractive area of the minimum. A full treatment of the stochasticity of the system lies beyond the scope of this paper and is a topic for future work. Nevertheless, from the work presented here it follows that stochastic effects occur at the pitch-angle minimum predicted by the drift approximation causing anomalous pitch-angle displacements.

The quantities of interest for the resulting pitch-angle development and isotropization of the electron distribution are the pitch-angles of the real electrons. The displacement is generally larger for smaller initial perpendicular velocities. This tendency is observed throughout the simulation. It reaches a maximum in the special case of vanishing $v_\perp$ which is not shown in Fig. 4.8. Electrons starting at $v_\perp = 0$ experience the strongest increase in pitch-angle and behave different from the other particles. However, zero perpendicular speed is a very special state for an electron and not relevant to the bulk reaction of the whole electron population.

The displacements in the mean of pitch-angles are depicted in Fig. 4.8. Three different regions can be identified: Below a value of about $v_\parallel \approx 10^8$ cm/s, the mean pitch-angle is shifted to lower values, in the intermediate range from $v_\parallel \approx 10^8 - 10^9$ cm/s there is a local maximum and above a value of $v_\parallel \times 10^{-9} \approx 10^9$ cm/s the pitch-angle displacement continuously increases with increasing $v_\parallel$. The latter one is interpreted as the increase by approaching resonances with the wave field starting at around $\sim 2 \cdot 10^{10}$ cm/s. Nevertheless, the interaction is non-resonant and belongs to the isotropization process to be investigated by this work. The result is a fan-like widening of the faster electrons with parallel velocities beyond the perpendicular thermal speed. At even higher initial velocity, this behavior dominates the isotropization process, ultimately yielding resonant pitch-angle scattering as observed in Messmer (2002).

The most interesting region from the point of view of non-resonant interaction is the region below $v_\parallel \approx 10^9$ cm/s. No resonances are present in this range and the response of the electrons is not expected to exhibit a strong dependence on the single particle kinetics. Nevertheless, Fig. 4.8 indicates a significant dependence on the initial parallel speed. Two regimes can be distinguished here: a negative shift below around $v_\parallel \approx 10^8$ cm/s and a positive shift above it. Although the negative shift at low $v_\parallel$ is of comparable magnitude to the positive shift in the intermediate range, it does not contribute much to the isotropization process, since at large angles a shift in pitch-angle does not result in much energy gain in perpendicular direction. In the intermediate range, stronger displacement in pitch-angle is observed and due to smaller, while still large, initial pitch-angles $\alpha_0$ more energy is transferred to the perpendicular direction.

The time scales for the displacement in pitch-angle are a direct result of the growth of the EF waves. Since the process is assumed to be adiabatic, at least in the ranges described in the above, the particle changes directly follow the growing waves. Hence, the characteristic time scale for the pitch-angle shift is given by the inverse of the linear growth-rate $\tau \sim 1/\gamma$ which is of order $\sim 1/\Omega_\nu$ in the case of EF waves.
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The simulations assume infinite coherent wave fields. In reality an electron will move into a region of space with entirely different wave phases within a time $\tau_c \approx \Delta z(\omega)/v_\parallel$, where $\Delta z(\omega)$ is the spatial dimension of coherent wave packets with frequency $\omega$. The size is determined by the formation of these waves, the EFI. The instability operates coherently on dimensions of $\Delta z(\omega) \ll f/\nabla_x f$, where $f$ is the electron distribution. For a scale length of $10^5$ cm in the acceleration region, $\Delta z$ may be roughly estimated to be $10^4$ cm and thus $\tau_c \approx 10^{-5}$ or a few proton gyro periods. This is consistent with the assumptions of the simulation of constant wave phases.

On the other hand, the above rough estimate suggests that an electron rapidly moves through regions with different wave phases. Every time it does so, scattering starts anew. The simulations show that in each step the pitch-angle increases a few degrees in the favorable regions of velocity space.

4.7 Summary

It has been established in a former work (Paesold & Benz, 1999) that the Electron Firehose instability (EFI) can arise as a result of electron acceleration in impulsive solar flares. Due to a velocity space anisotropy ($T_\perp < T_\parallel$) in the accelerated electron distribution, Electron Firehose (EF) waves are non-resonantly excited.

The kinetic response of the electrons in terms of pitch-angle development to the non-resonant instability is investigated in the work presented herein. Test particle simulations are carried out that are interpreted by a drift kinetic approach. The results clearly show that electrons are non-resonantly scattered to higher pitch-angles. It is caused by $E \times B$ drift, curvature drift and approximate conservation of magnetic moment.

At small wave field strengths the drift kinetic approach yields a good approximation to the system in both an electron in a single wave and in a spectrum of waves. By increasing the field strength in the single wave simulation, an upper limit for the validity of the drift approximation is found at a value of $e_k/B_0 \approx 10^{-4}$. In the case of a spectrum of waves the limit is not clearly defined. Due to variations in the magnitude of the total magnetic field $|B|$ adiabaticity can only be ensured for very small field strengths. At larger values of $e_k$ the electron orbits become stochastic and the resulting pitch-angle displacement cannot be described by the approximation.

In support of this conclusion it is found, that the particle response to the EF waves exhibits variations dependent on the particle velocities. Although the instability is hydrodynamic and does therefore not depend on single electron kinetics, a pronounced dependence of the electron response on the initial parallel velocity is observed. It is therefore expected that some intermediate velocities and very fast velocities isotropize more rapidly. Thus the evolution of the particle distribution in velocity space can develop peculiar shapes rather than maintaining an ellipsoidal form.

The simulations show that the isotropization needed in a solar particle acceleration context cannot be achieved in one step by a single set of waves. Nevertheless, the spatial inhomogeneity in the acceleration region causes a frequent change in the wave's phase.
relations resulting in multiple scattering, that is capable of isotropizing the electron distribution at a rate that is required for acceleration.
Chapter 5

Acceleration and Enrichment of $^3$He in Impulsive Solar Flares by Electron Firehose Waves

ABSTRACT: A new mechanism for acceleration and enrichment of $^3$He during impulsive solar flares is presented. Low-frequency electromagnetic plasma waves excited by the Electron Firehose Instability (EFI) can account for the acceleration of ions up to 1 MeV amu$^{-1}$ energies as a single stage process. The EFI arises as a direct consequence of the free energy stored in a temperature anisotropy ($T_{\parallel}^e > T_{\perp}^e$) of the bulk energized electron population during the acceleration process. In contrast to other mechanisms which require special plasma properties, the EFI is an intrinsic feature of the acceleration process of the bulk electrons. Being present as a side effect in the flaring plasma, these waves can account for the acceleration of $^3$He and $^4$He while selectively enhancing $^3$He due to the spectral energy density built up from linear growth. Linearized kinetic theory, analytic models and test particle simulations have been applied to investigate the ability of the waves to accelerate and fractionate. As waves grow in both directions parallel to the magnetic field, they can trap resonant ions. Plausible models have been found that can explain the observed energies, spectra and abundances of $^3$He and $^4$He.

5.1 Introduction

Solar flares are commonly divided into two different classes: impulsive and gradual (Cane, McGuire, & von Rosenvinge, 1986). The division into these two categories can be done on the basis of the duration of their soft X-ray emission (Pallavicini et al., 1977).
But it is not only the time scale of the events that justifies the distinction: the energetic particles observed in space from impulsive flares exhibit strong abundance enhancements over coronal values (Lin, 1987; Reames, 1990, and references therein). Impulsive flares are usually dominated by energetic electrons and are characterized by $^3\text{He}/^4\text{He}$ ratios at 1 MeV amu$^{-1}$ energies that are frequently 3 to 4 orders of magnitude larger than the corresponding value in the solar corona and solar wind where $^3\text{He}/^4\text{He} \sim 5 \cdot 10^{-4}$. They also exhibit enhanced $^4\text{He}/\text{H}$ and Fe/C ratios. Although the occurrence of $^3\text{He}$ and $^{56}\text{Fe}$ enrichments are correlated in impulsive flares, the ratio $^3\text{He}/\text{Fe}$ shows huge variations as observed by Mason, Dwyer, & Mazur (2000). This suggests, that different mechanisms are responsible for the acceleration of the two species. Gradual flares usually have large energetic proton fluxes, small $^4\text{He}/\text{H}$ and do not show large $^3\text{He}/^4\text{He}$ or Fe/C enhancements in the energetic particles, although approximately 5% admixture of supra-thermal remnant particles from impulsive flares have been observed in gradual events by Tylka et al. (2001). The standard interpretation for these observations is that the energetic particles in impulsive events origin in the energy release region on the sun while the energetic particles in gradual events are accelerated via shocks, either coronal or interplanetary (Lin, 1987; Luhn et al., 1987).

Abundance ratios therefore are a valuable diagnostics for the flaring plasma itself and in particular the specific acceleration mechanism for the energetic particles. The selectivity of the mechanism, especially for $^3\text{He}$ and $^4\text{He}$ indicates resonant processes such as gyro resonant interaction of plasma waves with the ions. Theoretical ideas therefore focus on the unique charge-to-mass ratio of $^3\text{He}$ which allows it to be selectively preheated or accelerated via gyro resonance.

A well-known theory of the initial set among theories for $^3\text{He}$ enhancement was published by Fisk (1978), explaining the preferential acceleration of the ions by electrostatic ion cyclotron (EIC) waves at a frequency in the vicinity of the gyrofrequency of $^3\text{He}$. The waves are excited by an electron current and interact with $^3\text{He}$ via cyclotron resonance. A large enhancement of $^4\text{He}/\text{H}$ is required in the ambient plasma for this instability to excite waves above the $^4\text{He}$ gyrofrequency.

More recently a theory was suggested by Temerin & Roth (1992) who accounted for the preferential $^3\text{He}$ acceleration proton electromagnetic ion cyclotron (H$^+$ EMIC) waves. These waves are driven unstable by non-relativistic (keV range) electron beams and their frequencies lie at around the $^3\text{He}$ gyrofrequency at almost perpendicular propagation. In Temerin & Roth (1992), auroral observations of keV electron beams and H$^+$ EMIC waves are taken as experimental evidence that H$^+$ EMIC waves also may acquire a substantial fraction (order of few percent) of the electron beam energy under coronal conditions. At the Sun, plasma emission from tens of keV electron beams on open magnetic field lines is the explanation for type III radio emission with its equivalent, the U-bursts on closed field lines (in loops) at beam energies of order keV. In order to excite H$^+$ EMIC waves fulfilling the requirements of this model, electron beams of much higher density than the observed ones have to be postulated (Miller & Viñas, 1993) with no direct observational evidence. Moreover, it is difficult to explain from a theoretical point of view that the free energy in the electron beam is transferred to the H$^+$ EMIC waves and not to the much faster growing ($\sim 4$ orders of magnitude in growth rate)
electron plasma waves.

In this work an alternative model for acceleration of $^3$He and its enhancement over $^4$He is presented. The approach is different from the models described above. While the models mentioned above postulate rather special plasma properties ($^4$He/H over-abundance in the pre-flaring plasma, dense low-energy beams) in order to produce the required plasma waves, the model presented herein explains the unique overabundance of $^3$He by plasma waves excited as an intrinsic feature of the electron acceleration process itself. Parallel propagating, left-hand polarized electromagnetic waves driven by the Electron Firehose instability (EFI) can account for $^3$He acceleration via gyro resonant interaction. As suggested in Paesold & Benz (1999) such Electron Firehose (EF) waves are excited by anisotropic electron distribution functions ($T_\parallel > T_\perp$) that occur in the course of the acceleration process of bulk electrons in solar flares.

Although the EF waves are not narrow banded around the gyrofrequency of $^3$He as the H$^+$ EMIC proposed by Temerin & Roth (1992), selectivity of the process is achieved by the natural profile of the spectral wave energy. No additional assumptions besides the electron anisotropy are needed and the model can be embedded as an intrinsic feature in acceleration scenarios such as transit-time damping, the currently most popular stochastic acceleration model.

In Sect. 5.2 the basics of the new acceleration scenario are described. The properties of the EF waves under coronal conditions are presented in Sect. 5.3.1. Analytical and numerical results on the heating rates are presented in Sect. 5.3.2 and Sect. 5.3.3. The mechanism for enhancement of $^3$He over $^4$He is described in Sect. 5.4, and heavier ions are discussed in Sect. 5.5. Section 5.6 concludes this work.

5.2 Basic Idea

Among the most promising scenarios for accelerating electrons to observed energies in impulsive solar flares is transit-time damping acceleration (Fisk, 1976; Stix, 1992), the magnetic analogon of Landau damping. The following scenario was presented by Miller et al. (1996): Electrons are accelerated from thermal to relativistic energies by resonance with low-amplitude fast-mode waves having a continuous broadband spectrum. The magnetic moment of the particle interacts with the parallel gradient of the magnetic field as in the well known Fermi acceleration (Fermi, 1949; Davis, 1956). Contrary to the classic Fermi process, transit-time damping involves small-amplitude magnetic compressions. While Fermi acceleration is the result of large numbers of particles being reflected by randomly moving magnetic compressions, transit-time damping is a process of rather resonant nature. In the limit of very small amplitudes, only particles of the same parallel speed as the wave phase velocity can be reflected by the magnetic compressions, i.e. $v_\parallel = \omega/k$ which is the Landau resonance condition. Interaction with a wave changes the particle's parallel speed and therefore allows it to interact with another wave of the continuous spectrum. In the average this process results in stochastic acceleration.

The wave spectrum is assumed to origin from cascading fast-mode waves, initially
excited at very long wavelengths as a direct output of the primary energy release in impulsive flares. The cascade channels the released energy through an inertial region to \( k \)-values small enough to accelerate electrons out of the thermal background population.

Other acceleration scenarios operating in impulsive solar flares have been proposed and can roughly be divided into three categories: Shock acceleration, acceleration by parallel electric fields and stochastic acceleration by MHD turbulence including the model described above. A detailed review can be found in Miller et al. (1997) and references therein.

Even though the nature of the actual acceleration mechanisms is unclear, all the above mechanisms that can account for accelerating the bulk of electrons up to the observed energies of \( \sim 20 \text{ keV} \) have an important feature in common: Particles are preferentially accelerated in parallel direction with respect to the background magnetic field. Parallel dc electric fields trivially accelerate only along the magnetic field while stochastic scenarios as transit-time damping act via small amplitude magnetic mirroring which is only capable of transferring energy in parallel direction if no additional scattering mechanism is provided (Lenters & Miller, 1998). Works of e.g. Wu (1984) and Leroy & Mangeney (1984) describe the parallel directed energization of electrons at the earths bow shock via shock drift acceleration with quasi perpendicular shocks.

The distribution function of the accelerated particle species in velocity space therefore is expected to become more and more anisotropic in the course of the energization process. While the perpendicular temperature remains virtually constant during the acceleration, the parallel velocity of the particles increases. If the energization in parallel direction is from a thermal level of some 0.1 keV up to 20 keV or more, the anisotropy in parallel direction is substantial. The plasma therefore is modeled by a bi-maxwellian with different temperatures in parallel and perpendicular direction. In doing so it is assumed that no efficient scattering mechanism is present that could maintain isotropy on acceleration time scales. Although it is not expected that acceleration retains a bi-maxwellian distribution function, it is a good approximation in order to describe propagation properties of plasma waves. As shown in Paesold & Benz (1999) the EFI can occur in such a situation and give rise to EF waves. The parallel EF waves are purely transverse electromagnetic and left-hand circularly polarized. The sense of polarization and can change in certain \( k \) ranges if a similar anisotropy of the protons is assumed. In the following the EF waves are left-hand polarized if not mentioned otherwise. A representative dispersion plot is depicted in Fig. 5.1.

Taking into account the uncertainties in the acceleration region, including possible pre-heating mechanisms, reasonable pre-flaring plasma conditions of an impulsive flare range within \( B_0 \approx 100 - 500 \text{ G} \) for the background magnetic field, \( n_\alpha \approx 10^9 - 10^{11} \text{ cm}^{-3} \) in number density and about \( T_\alpha \approx 10^6 - 10^7 \text{ K} \) for the proton and electron temperature (Pallavicini et al., 1977). For the numerical example herein the following pre-flaring plasma parameters are chosen: \( n_e = n_p = 5 \cdot 10^{10} \text{ cm}^{-3}, T_{\perp,\parallel}^e = T_{\perp,\parallel}^p = 1 \cdot 10^7 \text{ K}, B_0 = 100 \text{ G} \).

Under these exemplary conditions the EF waves propagate below about \( 3 \cdot \Omega_\parallel \). Several ion gyrofrequencies are indicated in Fig. 5.1 and they lie well within the wave spectrum. Resonant acceleration occurs if the condition \( \omega - k_{\parallel}v_{\parallel} - i\Omega/\gamma = 0 \) is satisfied.
Here, \( v_\parallel \) and \( \gamma \) are the parallel particle speed and Lorentz factor, and \( \Omega \) is the gyro-frequency of the according particle. This condition is well satisfied for ions like \( ^3\text{He}^{2+} \) and \( ^4\text{He}^{2+} \). The model presented herein assumes acceleration of the ions via the most effective gyro resonance at \( l = 1 \), the so called cyclotron resonance.

Due to the symmetry of the distribution function with respect to \( v_\parallel \), the dispersion of the EF waves is the same for the transition from \( k \rightarrow -k \) while keeping left-hand polarization and \( \omega > 0 \) (Hollweg & Völk, 1970). The resulting wave field from the EFI therefore consists of waves propagating parallel to the magnetic field (positive k-branch) as well as anti-parallel (negative k-branch). This property of the EF waves is of special interest and crucial for efficient acceleration: A unidirectional wave field of electromagnetic waves always exhibits a force parallel to the background magnetic field pushing the particle out of resonance, limiting the acceleration time and, hence, the reachable energies. While other models (e.g. Roth & Temerin 1997) solve the problem by imposing a background magnetic field geometry, i.e. a field gradient, to force the particle to regain resonance, this is not necessary for the case of the EF waves: resonant acceleration by counterpropagating electromagnetic waves leads to oscillating parallel forces and does not drive the particle permanently out of resonance. In contrary they naturally 'trap' the particle in the wave fields by bouncing it from resonance with the positive branch to resonance with the negative branch of the wave field and vice versa. This ping-pong trapping strongly enhances the total resonant interaction time of the particle with the wave and allows significant acceleration in perpendicular velocity. The mechanism works even for small wave fields where 'normal' wave-particle trapping is not efficient. Similar situations have been investigated by e.g. inferring counterpropagating Alfvén waves to the problem of solar flare proton acceleration (Barbosa, 1979) and cosmic-ray acceleration (Skilling, 1975). This acceleration can be regarded as a special case of wave turbulent stochastic acceleration with waves explicitly counterpropagating in one direction.

The differences in enrichment of \(^3\text{He}\) and \(^4\text{He}\) is the result of varying growth of the waves at different frequencies. As the waves grow from a thermal level to a saturated state, the spectral energy density develops according to the growth rates from linear theory depicted in Fig. 5.1. Since the waves are excited non-resonantly by the electrons, the saturated wave energy spectrum can be approximated from the profile of the linear growth rate. This is not true for resonantly excited instabilities, where wave modes can still grow while others have already saturated. If each mode is in resonance with only the particles fulfilling the resonance condition, some modes can grow longer than others and the saturated energy spectrum of the waves has to be determined by other methods. However, if the instability is non-resonant as the EFI, each mode grows according to the total free energy available (the driver is basically a pressure anisotropy of the electrons) and the whole electron population contributes to the growth of each mode. All modes therefore saturate at the same time, namely when the pressure anisotropy drops below the threshold, and the spectral energy is frozen at that time. The wave dispersion is altered by the decreasing anisotropy as result of the erosion of the particle distribution. However, a former analysis shows (Paesold & Benz, 1999) that this does not severely alter the slope in the growth rate around \(^3\text{He}\) an \(^4\text{He}\).
Although the differences in growth rates are quite small in the region of interest (see Fig. 5.2), the differences in wave energy will become significant after several growth times $\tau = 1/\gamma$. As a rough estimate one obtains, by assuming constant growth rates in time, that two modes with initial growth rates $\gamma_A = \gamma_{\text{Max}}$ and $\gamma_B = 0.9 \, \gamma_{\text{Max}}$ and equal energy content have a ratio in energy of $W_B/W_A \sim 0.1$ after a time $\tau = \gamma_A t$ of only $\approx 10$. This estimate illustrates the rather strong influence of growth rate on the resulting wave energy level and, hence, on the selectivity of cyclotron acceleration due to the $\gamma$ profile in wave frequency. The analysis in Paesold & Benz (1999) shows that the positive slope in the growth rate profile around the $^3\text{He}$ and $^4\text{He}$ gyrofrequency that is needed to selectively enhance $^3\text{He}$ is a stable feature of the EFI and not very sensitive to changes in the plasma parameters for solar pre-flaring conditions.

In the following it will be established that the mechanism meets the following requirements for the observed $^3\text{He}$ enrichment:

(i) The accelerator has to energize the ions above $\sim 1\,\text{MeV amu}^{-1}$ on time scales of 1 s.

(ii) $^3\text{He}/^4\text{He} > 0.1$ above energies of about 1 MeV amu$^{-1}$.

(iii) A time integrated total of about $10^{31}$ $^3\text{He}$ nuclei is required to account for the observed particle fluxes (Reames et al., 1994).

(iv) The $^3\text{He}$ spectrum is harder than the $^4\text{He}$ spectrum in the range of $0.4 - 4.0\,\text{MeV amu}^{-1}$ (Moebius et al., 1982).

(v) $^3\text{He}$ exhibits a turnover in the particle energy spectrum at around a few $\sim 100\,\text{keV amu}^{-1}$ (Mason et al., 2000).

5.3 Acceleration

5.3.1 Electron Firehose Wave Properties

Parker (1958) pointed out that a magnetized plasma with a pressure anisotropy in parallel direction to the magnetic field can become unstable to low frequency Alfvén waves. This instability is known as the Firehose instability and is of completely non-resonant nature: neither the electrons nor the protons are in resonance with the excited waves. An extension of this instability to higher frequencies was presented by Hollweg & Völk (1970) and Pilipp & Völk (1971). This branch of the instability has been termed the Electron Firehose Instability. Here, the bulk of the protons is resonant with the waves which are non-resonantly excited by the electrons. The electrons are anisotropic and drive the waves while the protons carry the wave. A typical dispersion relation of the EF waves is displayed in Fig. 5.1. The dispersion of the EF waves is computed by IDLWhamp (Paesold, 2002), an easy to use IDL interface to the WHAMP code originally developed by Rönnmark (1982). The code provides the user with the full solution of the dispersion equation in linearized kinetic theory.
Figure 5.1: Dispersion relation of Electron Firehose waves calculated from linear theory. The frequency $\omega_r$ (solid curve) and growth rate $\gamma$ (dashed curve) of the EF waves are shown normalized to the proton gyrofrequency as function of parallel wavenumber $k_\parallel$ normalized to the inertial length of the protons. The plasma parameters are $T_{e//}/T_{e\perp} = 15$, $T_{p//}/T_{p\perp} = 2$, $T_{p//} = T_{p\perp} = 1 \cdot 10^7$ K, $n_e = n_p = 5 \cdot 10^{10}$ cm$^{-3}$ and $B_0 = 100$ G. The sign convention is that $k > 0$ and $\omega > 0$ mean left-hand circularly polarized; $\gamma > 0$ refers to growing modes. The dotted horizontal lines indicate the gyrofrequency of the according ion species.

The influence of the presence of other majority ions such as $^4$He on the EF wave dispersion is negligible. For a plasma consisting of 5% $^4$He and 95% H the wavenumber of maximum growth and the frequency at maximum growth is shifted by values of order of 1% with respect to the values of a pure H plasma. Ions other than H$^+$ have therefore been omitted in the following when computing the dispersion relation of the EF waves. Due to the assumed anisotropy of the protons the mode is right-hand polarized ($\omega_r < 0$) at small values of $k$ in Fig. 5.1. According to Hollweg & Völk (1970) the condition for the change of polarization sense to occur is an additional proton anisotropy $T_{p//} / T_{p\perp} \geq 2$. Such a proton anisotropy can result e.g. from transit-time acceleration as described for the electrons in Sect. 5.2. Note, that the proton anisotropy is not a necessary condition for instability. It has been taken into account herein only for the sake of generality. The dependence of the maximum growth rates on proton and electron anisotropy can be seen in Fig. 5.3. The proton anisotropy has been varied from 1 to larger values for three values of the electron anisotropy. Only left-hand modes have been taken into account. With increasing proton anisotropy the frequency at maximum growth approaches and finally crosses the $^3$He gyrofrequency. With decreasing electron anisotropy less anisotropic protons are needed to shift the most growing mode to the $^3$He gyrofrequency. This shows that for reasonable anisotropies the slope in the growth rate profile in the vicinity
of the gyrofrequency of \(^3\)He is always positive, such that waves around the \(^3\)He resonance grow faster than waves around the \(^4\)He gyrofrequency. This is a very stable feature of the EF wave spectrum and can be expected, once above the instability threshold, for a large range of plasma parameters that allow instability and reasonable anisotropies of protons and electrons (for a detailed analysis of EFI thresholds and spectra in dependence on plasma parameters see Paesold & Benz, 1999).

An additional and faster growing mode can be found at oblique propagation angles (Paesold & Benz, 1999; Li & Habbal, 2000). However, this mode is not of interest for ion acceleration. The oblique mode exhibits growth rates at maximum of about two orders of magnitude larger than the parallel mode (Li & Habbal, 2000). Since the growth rate lies then well above the ion gyrofrequencies, the according ions are unmagnetized and cannot damp the waves, hence cannot gain energy (Gary, 1993). At larger angles, the oblique mode is purely growing and therefore is not capable of accelerating particles via cyclotron resonance. Although the oblique mode can reduce anisotropy and grows faster than the parallel mode, it is a reasonable assumption, that there is enough energy available to drive the parallel mode.

### 5.3.2 Heating Rates

#### Unidirectional Propagation

The ion’s energy change is dominated by the resonant interaction with the wave’s electric field, which is decomposed into its parallel component \( E_\parallel = \hat{E}_\parallel \cos \Psi \), where \( \Psi = \)
Figure 5.3: Maximum growth rate $\gamma_{\text{Max}}/\Omega_H$ vs. frequency $\omega/\Omega_H$. For each of the three lines the electron anisotropy is constant and the proton anisotropy is varied. The most upper line (triangles) depicts $T_{H}/T_{H} = 20$, the middle line (diamonds) $T_{H}/T_{H} = 15$ and the bottom line (asterisks) $T_{H}/T_{H} = 14$. The small numbers refer to the according values of $T_{H}/T_{H}$. The vertical dotted line indicates the $^3\text{He}^{2+}$ gyrofrequency.

$k_{||} x + k_1 x - \omega t$ is the wave’s phase at the location of the ion, and two circularly polarized components $E_{\pm} = E_{\pm}(\cos \Psi, \mp \sin \Psi)$ with amplitudes $E_{\pm} = (E_x \pm E_y)/2$. In the Fourier domain we decompose the wave into a set of monochromatic waves separated by a frequency interval $\Delta \omega$. In one of these monochromatic waves with amplitudes $\hat{E}_{\parallel}$ and $\hat{E}_{\pm}$ and frequency $\omega$, the ion gains or loses energy according to

$$W = qv_{\perp} \hat{E}_{\pm} \cos \phi_{\pm} + qv_{\parallel} \hat{E}_{\parallel} \cos \Psi,$$

where the relative angle $\phi_{\pm} = \theta \pm \Psi$ between the wave electric field and the perpendicular velocity vector $(v_x, v_y) = v_{\perp}(\cos \theta, \sin \theta)$ of the ion, moving with the instantaneous gyro phase $\theta(t)$.

The ion may be driven into or out of cyclotron resonance due to changes in the ion’s parallel speed, which in turn depends on the wave’s electric and magnetic field. The latter is obtained via Faraday’s law $\mathbf{B} = -\nabla \times \mathbf{E}$ as $(1/\omega) [\pm (k_{||} \hat{E}_{\pm}) \sin \Psi, (k_{||} \hat{E}_{\pm} - k_{\perp} \hat{E}_{||}) \cos \Psi, \mp (k_{\perp} \hat{E}_{\pm}) \sin \Psi]$, so that the parallel acceleration is

$$m \ddot{v}_{||} = q \left[ \frac{k_{||} v_{\perp}}{\omega} \hat{E}_{\pm} \cos \phi_{\pm} + \left(1 - \frac{k_{\perp} v_{\perp} \cos \theta}{\omega} \right) \hat{E}_{||} \cos \Psi \right].$$

A possible mirror force due to a background magnetic field gradient has been omitted since $B_0$ is assumed to be homogeneous herein.
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The instantaneous ion angular frequency, $\dot{\varphi}_i$, is given by the ratio of the force on the ion, perpendicular to the velocity and the background magnetic field, and the moment perpendicular to the magnetic field. Considering only one monochromatic wave at frequency $\omega$, the phase equation $\dot{\phi}_\pm = \theta \pm \Psi$ is given by

$$\dot{\phi}_\pm \approx -\Omega_i \pm (k_{||}v_{||} - \omega + k_{\perp}v_{\perp}\cos \theta) - \frac{q\hat{\epsilon}_\pm}{mv_{\perp}} \left(1 - \frac{k_{||}v_{||}}{\omega}\right) \sin \phi_\pm,$$  \hspace{1cm} (5.3)

where $\Omega_i$ is the gyrofrequency of the ion $i$.

In case of the EF waves, where $k_{\perp} = 0$, $\hat{\epsilon}_{||} = 0$, and $\hat{\epsilon}_+ = 0$, Eqs. (5.1) and (5.2) simplify to

$$\dot{W} = q\hat{\epsilon}_- v_{\perp} \cos \phi_-, \hspace{1cm} (5.4)$$

$$\dot{v}_{\perp} = \frac{q\hat{\epsilon}_-}{m} \cos \phi_- \left(1 - \frac{k_{||}v_{||}}{\omega}\right), \hspace{1cm} (5.5)$$

$$\dot{v}_{||} = \frac{q\hat{\epsilon}_- k_{||}v_{||}}{m} \cos \phi_-.$$

For a particle $i$ in resonance with a monochromatic EF wave the frequency mismatch parameter $\xi = \Omega_i + k_{||}v_{||} - \omega$ vanishes. The right-hand side of Eq. (5.3) therefore reduces to a nonlinear pendulum equation for the phase $\phi_-$. In resonance $\phi_-(t)$ can have a stable solution when it stays close to 0. In this case the particle is accelerated in parallel direction until it is expelled from resonance (Equation 5.6). For particles not in resonance, $\phi_-$ changes rapidly and no energy is gained in the time average.

In a set of $N$ monochromatic waves, numbered by $j = 1, \ldots, N$ with equal amplitudes $\hat{\epsilon}_{-j} = \sqrt{-E_{rms}/N}$ and equally spaced by a frequency $\Delta \omega$ with $\delta \omega = N\Delta \omega$, the ion's perpendicular motion has a phase $\phi = \theta - \Psi_j$ with respect to each mode. This gives a set of $N$ phase equations

$$\dot{\phi}_j = -\Omega_i - k_{||}v_{||} - \omega_j - \sum_{l=1}^{N} \frac{qE_{rms}}{mv_{\perp} \sqrt{N}} \left(1 - \frac{k_{||}v_{||}}{\omega_l}\right) \sin \phi_l,$$  \hspace{1cm} (5.7)

and Eqs. (5.5), (5.6) have to be summed over the contributions of the $N$ modes.

The equations of motion are in general not integrable because of the $\sin \phi_l$ terms in Eq. (5.7). However, a maximum heating rate and the typical time after which an ion is driven out of resonance by the mean parallel force of the uni-directional propagating waves are estimated. The latter limits the total energy gain of the ion.

Within a time interval $\tau$, the energy gain of the ion is mainly due to the wave modes in a bandwidth $2\pi/\tau$ around the frequency $\omega \approx \Omega_i + k_{||}v_{||}$. The wave may be described by a mode spacing $\Delta \omega = 2\pi/\tau$, yielding $\hat{\epsilon}_- \approx E_{rms}/\sqrt{\delta \omega \tau/2\pi}$. At best, the ion stays in resonance with one mode $j_0$ such that $\phi_{j_0} \approx 0$ over the whole time period $\tau$. This gives, using Eq. (5.5),

$$v_{\perp}(t + \tau) - v_{\perp}(t) \approx \sqrt{2\pi} \frac{qE_{rms}\tau}{m\sqrt{\delta \omega \tau}}, \hspace{1cm} (5.8)$$
and if \( \tau \) is sufficiently large so that \( v_\perp(t) \ll v_\perp(t + \tau) \), the heating rate per mass is at most

\[
v_\perp \dot{v}_\perp \approx \frac{2\pi q^2 E_{\text{rms}}^2}{m^2 \delta \omega}
\]

\[
\frac{W_\perp}{m} \approx 6.1 \left[ \text{Mev} \frac{\text{amu s}}{100 \text{ V m}^{-1}} \right] \left( \frac{E_{\text{rms}}}{10^6 \text{ rad s}^{-1}} \right) \left( \frac{Q}{\text{A}} \right)^2 \tag{5.9}
\]

independent of the length of the time interval \( \tau \). Thus, Eq. (5.9) predicts a linear increase of energy in time. The ratio \( Q/A \) is the charge-to-mass ratio of the ion in atomic units.

The resonant heating is limited to the time, during which the ion is in resonance with one of the wave modes. The ion dynamics are dominated by the time scales for parallel and perpendicular acceleration

\[
\tau_\parallel = \frac{v_\parallel}{\dot{v}_\parallel} = \frac{mv_\parallel \omega}{q \hat{\epsilon} - k_\parallel v_\perp}
\]

\[
\tau_\perp = \frac{v_\perp}{\dot{v}_\perp} = \frac{mv_\perp}{q \hat{\epsilon} - 1 - k_\parallel v_\parallel / \omega} \tag{5.10}
\]

The initial conditions are usually \( v_{\perp0} \approx v_\parallel \) with \( k_\parallel v_\parallel \ll \delta \omega \) and \( \tau_{\perp0}/\tau_{\parallel0} < 1 \). With increasing \( v_\perp \), the time scale \( \tau_\perp \) increases and the time scale \( \tau_\parallel \) decreases. At time \( t \approx t_r \), the ion eventually reaches a parallel speed \( |v_\parallel| \) of order \( \delta \omega/k_\parallel \) and is driven out of resonance. The acceleration time scale at \( t_r \) is

\[
\tau_{\parallel, r} \approx \frac{m \delta \omega \sqrt{\delta \omega \tau_{\parallel, r}}}{\sqrt{2\pi q E_{\text{rms}} k_\parallel} v_\perp(t_r)} \approx \frac{m^2 \omega^2 \delta \omega^3}{2\pi q^2 E_{\text{rms}}^2 k_\parallel^4 v_\perp^2(t_r)} \approx \frac{m^4 \omega^2 \delta \omega^4}{(2\pi)^2 q^4 E_{\text{rms}}^4 k_\parallel^4 t_r} \tag{5.12}
\]

Estimating \( t_r \) shows that the total time of the resonant acceleration process is not much longer than \( \tau_{\parallel, r} \), and thus from Eq. (5.13) follows

\[
t_r \approx \frac{m^2 \omega \delta \omega^2}{2\pi q^2 E_{\text{rms}}^2 k_\parallel^3} \approx 1.6 \frac{10^3}{\omega} \left( \frac{\delta \omega}{\omega} \right)^2 \left( \frac{A}{Q} \right)^2 \left( \frac{100 \text{ V m}^{-1}}{E_{\text{rms}}} \right)^2 \left( \frac{1 \text{m}^{-1}}{k_\parallel} \right) \left( \frac{B_0}{0.01 \text{ T}} \right)^4 \tag{5.13}
\]

With Eq. (5.9) the corresponding energy gain is \( v_\perp^2 \approx \omega \delta \omega / k_\parallel^2 \), which is only a few keV amu\(^{-1}\) for the given parameters. It is too little to explain the observation, assuming realistic conditions.
Bidirectional Propagation

Further acceleration is possible in the situation of the EFI where counterpropagating electromagnetic waves of the same sense of polarization are present. The basic idea is, that the time average of the wave forces in parallel direction of counterpropagating waves averages out to some extent. That way, the ion can stay in resonance for much longer time.

In the following left-hand circular polarization is assumed. Plus and minus signs in index positions therefore indicate the direction of propagation and no longer sense of polarization. The parallel acceleration now reads

\[ \dot{v}_\parallel = \sum_{j=1}^{N} \frac{q k_{j} v_{\perp}}{\omega_{j} m} \left[ e_{j} \cos \Phi_{+,j} - e_{-j} \cos \Phi_{-,j} \right], \]

where the phases \( \Phi_{\pm,j} = \theta \mp k_{j} v_{\parallel} t + \omega_{j} t \) are introduced. Then \( \Phi_{-,j} \) can be expressed by \( \Phi_{+,j} \) and we rewrite Eq. (5.15) as

\[ \dot{v}_\parallel = \sum_{j=1}^{N} \frac{q k_{j} v_{\perp}}{\omega_{j} m} \left[ (e_{j} - e_{-j}) \cos \Phi_{+,j} \right. \]

\[ \left. + 2 e_{-j} \sin (\theta + \omega_{j} t) \sin \left( k_{j} v_{\parallel} t \right) \right]. \]

The first term in the bracket is similar to the unidirectional case (Equation 5.6) except that the difference between the bidirectional fields enters now and reduces the parallel acceleration. As in the unidirectional case this term is nearly constant in resonance, i.e. \( \Phi_{+,j} \approx 0 \). The second term in the bracket contains a product of two sinus terms. At times \( t_v < 2\pi/k_{j} v_{\parallel} \) there is a mean change in \( v_{\parallel} \) when averaging over \( t \) due to the first sinus term. The second sinus term, \( \sin \left( k_{j} v_{\parallel} t \right) \) leads to an oscillatory motion in \( v_{\parallel} \) for \( t_v > 2\pi/k_{j} v_{\parallel} \) and thus slows down the escape from the resonance region and, eventually, allows the regain of resonance. The maximum energy, that can be reached by the ion therefore can be increased. The dynamics of \( v_{\parallel} \) (Equation 5.16) is nonlinear and treated in numerical simulations with the exact wave spectrum and dispersion relation of the counterpropagating waves.

5.3.3 Test Particle Simulation

The relativistic equations of motion for a particle of charge \( q \) and mass \( m \) in a field of \( N \) waves and a homogeneous background magnetic field \( B_0 \) are given by

\[ \frac{d\mathbf{p}}{dt} = q \frac{\mathbf{v}}{c} \times \mathbf{B}_0 + q \sum_{k=1}^{N} \left( \mathbf{E}_k + \frac{\mathbf{v}}{c} \times \mathbf{B}_k \right) \]

\[ \frac{d\mathbf{x}}{dt} = \mathbf{v}, \]

where \( \mathbf{x} \) is the particle position vector, \( \mathbf{v} \) the velocity and \( \mathbf{E}_k \) and \( \mathbf{B}_k \) are the electric and magnetic fields of the wave \( k \).
Test particle trajectories are calculated by integrating a dimensionless form of Eqs (5.17) and (5.18) with a standard leap-frog mover following Birdsall & Langdon (1991). The time integration scheme is displayed in Fig. 5.5. The system is prepared in a state where $\vec{x}$, $\vec{E}$ and $\vec{B}$ is known at an integer multiple of the time step $\Delta t$. This procedure deviates a little bit from the one in Birdsall & Langdon (1991), since it is a test particle code and the magnetic and electric fields can be computed at every time and are not the result of a current or charge density. There is no need to interpolate $\vec{B}$. By using $\vec{E}_t$, $\vec{B}_t$, and $\vec{x}$ and a discretized version of Eq. (5.17), the velocity $\vec{v}_{t-\Delta t/2}$ can be advanced to $\vec{v}_{t+\Delta t/2}$. By knowing $\vec{v}_{t+\Delta t/2}$, the particles space coordinate can now be pushed to $\vec{x}_{t+1}$ with the aid of a discretized version of Eq. (5.18). The fields now can be updated to $\vec{E}_{t+1}$ and $\vec{B}_{t+1}$ and the system is in the same iterational state as at the beginning. The new values overwrite the old values and none are retained more than one time. The code used herein is a leap-frog version of the code used by Miller & Viñas.
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At $n^{1/2}$

Figure 5.5: Temporal layout of field and particle quantities used in the leap frog integration of Eqs. (5.17) and (5.18).

(1993) and has been tested against their version. The simulation results of both codes are in excellent agreement on time scales relevant for the analysis herein.

The time step in the simulation was $\tau = 0.1\Omega_v^{-1}$. The background magnetic field is considered to be homogeneous. The wave frequencies, parallel wavenumbers and relative strength of electric field components are obtained from the IDLWhamp code described in Sect. 5.3.1. The electric field strength $E_k$ of one single wave $k$ is obtained from growth according to linear theory and the magnetic field components are calculated according to Faraday’s law. A wave field consisting of 2000 monochromatic waves has been applied, confined to a range of frequencies where the growth rate does not drop below $\sim 60\%$ of the maximum growth rate. All properties of the wave field besides the relative phases are symmetric with respect to the zero of $k_\parallel$. The phases were chosen randomly for all waves independent of the direction of propagation. Only waves propagating parallel to the background magnetic field are considered. Some properties of the applied wave spectrum are displayed in Fig. 5.4. The waves are purely transverse and left-hand circularly polarized.

The spectral energy density of the wave field is obtained by letting the waves exponentially grow according to $W(\tau, k) = W_0(k) \exp(2\gamma(k) \cdot \tau)$ from a thermal level $W_0(k) = k_BT$ up to an energy density of $W(\tau, k)$ at a time $\tau$ and $\gamma(k)$ is assumed to be constant during this time. The time of growth has been chosen such, that the total energy density in the wave field is $W_{tot}(\tau)/U_0 = 0.01$ at a time $\tau$, where $U_0 = B_0^2/8\pi$ is the energy density of the background magnetic field.

The plasma parameters are $n_e = n_p = 5 \times 10^{10}$ cm$^{-3}$, $T_\perp^e = T_\perp^p = 1 \times 10^7$ K, $T_\parallel^e/T_\perp^e = 15$, $T_\parallel^p/T_\perp^p = 2$ (see Fig. 5.1) and the background magnetic field is $B_0 = 100$ G.

To reduce computing time, the simulations have been divided in two parts: I) For each of the two ion species, a small number of particles have been followed for a long time ($\sim 1$ s) to establish the ability of the process to accelerate ions to the requested energies (1 MeV amu$^{-1}$). About 500 particles were computed in order to determine also the spread in energy of the particle population. II) A larger number of ions (5000 particles for each species) have been followed for shorter times ($\sim 0.052$ s) to establish...
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Figure 5.6: Temporal history of a typical $^3$He$^{2+}$ test particle. (a) The kinetic energy in keV amu$^{-1}$ of an initially thermal $^3$He ion (100 eV amu$^{-1}$) vs. time in units of $10^4\Omega_H^{-1}$. (b) Distance traveled by the particle along the background magnetic field $B_0$. (c) Perpendicular and (d) parallel momentum of the particle vs. time.

the energy distribution which then was applied to the results of the first runs.

A representative run of the part \textit{I)} is depicted in Fig. 5.6. The initial energy of the ion is 100 eV amu$^{-1}$, corresponding to a thermal velocity at a temperature of $1 \cdot 10^7$ K, and the pitch-angle cosine is $\mu = 0.5$. The heating rates in the simulation are not very sensitive to variations in the initial energy of the ion. All particles start at space coordinates equal to zero. For each particle the phases of the wave field have been chosen randomly. The total energy in the wave field is $U_W/U_0 \approx 0.01$ corresponding to the wave electric field amplitudes $E_k/B_0$ depicted in Fig. 5.4. The plasma parameters are chosen as described in Sect. 5.2.

The typical energy evolution of a $^3$He ion shows intervals of zero gain (Figure 5.6a). These are times when the ion looses resonance with the wave field due to a large excursion of $v_\parallel$ from its oscillatory behavior described by Eq. (5.16). During these intervals, the ion oscillates in the wave field without being significantly accelerated. Due to non-resonant interaction with the waves, the ion is scattered back into resonance and can be further accelerated. The loss and gain of resonance is mirrored in panel (b) and (d) of Fig. 5.6. During the intervals of no resonance (e.g. $(t \times \Omega_H) \times 10^{-4} \approx 8 - 15$), the ion propagates
freely and the coordinate along $B_0$ (Figure 5.6b) is just a straight line, increasing or decreasing proportional with time. As can be seen in Fig. 5.6d the parallel momentum no longer oscillates around zero in the above time interval. When the ion re-enters resonance, the parallel momentum starts oscillating around zero, indicating ping-pong trapping of the particle, and the coordinate along $B_0$ does not linearly change anymore.

To further illustrate the ping-pong behavior, the so-called frequency mismatch parameter $\xi = \omega_r - k v_{||} - \Omega_3$ (where $\Omega_3$ is the $^3$He$^{2+}$ gyrofrequency) has been plotted in Fig. 5.7. Resonant interaction of the particle with a wave means $\xi \approx 0$, reproducing the resonance condition. Panel (a) of Fig. 5.7 displays a small portion of the time serie of Fig. 5.6a and the three vertical lines indicate times at which $\xi$ was computed. As can be seen in Fig. 5.7b, the particle is out of resonance at times $t_1$ and $t_3$ while it is resonantly interacting with both wave branches at time $t_2$. Accordingly, the parallel normalized momentum $p_x/mc$ does not exhibit any oscillatory behavior around zero at times $t_1$ and $t_3$ in Fig. 5.7a but clearly oscillates around zero at $t_2$.

Acceleration is clearly very efficient with an averaged systematic energy gain $(dE/dt)$ of $\approx 1.2$ MeV amu$^{-1}$ s$^{-1}$ (Figure 5.6a). A linear energy increase is predicted by Eq. (5.9). The paths of more than 500 $^3$He ions have been simulated and exhibit similar energization.

5.4 Abundance Enhancements

Among the over-abundant ions observed during impulsive solar flares, $^3$He takes a special position. Whereas most ions exhibit an enhancement by factors of 1-10 with respect to coronal values, $^3$He shows an excess of $\sim 2000$. Since it has been established that the selective enhancement of ion abundances correlates with the charge-to-mass ratio (Reames, 1998) and hence the gyrofrequency of the according ions, resonant wave-particle interaction is generally believed to be the accelerating mechanism. This led earlier theories to the approach of finding special plasma waves existing only in a narrow range in the vicinity of the $^3$He gyrofrequency. In order to generate the appropriate waves (e.g. EIC waves (Fisk, 1978), H$^+$ EMIC waves in Temerin & Roth (1992)) peculiar plasma conditions have to be postulated. In the first case a pre-flare enhancement of $^4$He over H and in the latter case dense low-energy beams have to be assumed in order to excite the waves in the appropriate frequency ranges.

A different approach is investigated in the work presented herein. EF waves, excited by an anisotropic velocity distribution function resulting from bulk acceleration of electrons, accelerate $^3$He and $^4$He. Selectivity is achieved by an energy density profile of the waves self-consistently produced from linear growth. For non-resonant instabilities, this profile is approximately true not only during the growth phase but also for the saturated state of the instability (see Sect. 5.2). As can be seen from Fig. 5.4 the resulting energy density profile clearly supports an enhanced acceleration of $^3$He above $^4$He. The energy density in the wave field is higher at the gyrofrequency of $^3$He than $^4$He. More energy is therefore available for the acceleration of $^3$He. The quantitative heating rates of an ensemble of ions of the two species have been determined by numerical simulations.
where the total energy density in the wave field has been treated as a free parameter.

In order to compare the simulation results to the observational values of $^3\text{He}/^4\text{He}$, abundances have to be compared at the same energies. Observations from the ISEE 3 spacecraft were taken in the 1.3–1.6 MeV amu$^{-1}$ channel (Reames et al., 1994) and observations by Moebius et al. (1982) with the ISEE 1 and ISEE 3 spacecraft were taken in the 0.4–4.0 MeV amu$^{-1}$ range. According to these measurements, a viable mechanism for $^3\text{He}$ acceleration and its enhancement above $^4\text{He}$ has to meet the requirements of Sect. 5.2: (i) Ion energies $\sim$ 1 MeV amu$^{-1}$ on time scales of 1 s, (ii) $^3\text{He}/^4\text{He} > 0.1$ above 1 MeV amu$^{-1}$, (iii) a total of about $10^{31}$ $^3\text{He}$ nuclei above 1 MeV amu$^{-1}$, (iv) $^3\text{He}$ spectrum harder than the $^4\text{He}$ spectrum in the range of 0.4–4.0 MeV amu$^{-1}$ and (v) a turnover in the particle energy spectrum of $^3\text{He}$ at around $\sim$ 100 keV amu$^{-1}$.

In the following, the presented model is investigated in view of these requirements.
Chapter 5. Acceleration and Enrichment of $^3$He by EF waves

Figure 5.8: Temporal evolution of thermal energy $E_{th}(t)$ as defined in Eq. (5.19). A maxwellian distribution has been fitted to an ensemble of 500 particles for each species.

The ability of the EF waves to accelerate the ions to the relevant energies has been shown in Sect. 5.3 and requirement (i) therefore is fulfilled. In order to address items (ii) - (v) the collective behavior of an ensemble of ions has to be investigated.

From part II) of the numerical simulations (5000 particles, short times) it is found that the population of ions in energy develops according to a maxwellian distribution function for a system of three degrees of freedom

$$f_\alpha(E, t) = \frac{n_\alpha}{\sqrt{2\pi E_{th,\alpha}^3(t)}} \sqrt{E} \exp\left(-\frac{E}{2E_{th,\alpha}^3(t)}\right),$$  \hspace{1cm} (5.19)

where $f_\alpha(E, t)$ has been normalized to $n_\alpha$, the density of species $\alpha$.

By fitting this distribution model to the runs of part I) (500 particles, long times) the thermal energy in time $E_{th,\alpha}^3(t) = 1/2 m_\alpha (v_{th,\alpha}^3(t))^2$, and hence the heating rates, were extracted. The function $E_{th,\alpha}^3(t)$ for $^3$He and $^4$He is depicted in Fig. 5.8. Clearly, $^3$He is accelerated faster than $^4$He and this causes an enhancement at a given energy (say 1 MeV amu$^{-1}$).

The ratio $^3$He/$^4$He above 1 MeV amu$^{-1}$ can now be calculated in dependence on acceleration time. We integrate

$$n_\alpha(t) = \int_{1 \text{ MeV amu}^{-1}}^{\infty} f_\alpha(E, t) \, dE,$$ \hspace{1cm} (5.20)

where $\alpha = 3, 4$ corresponds to the according He isotope, and form the ratio $(^3$He/$^4$He)$\times(t) = n_3(t)/n_4(t)$. It is found that the required abundance ratio of $^3$He/$^4$He $\sim$
0.1–1 is reached in the time range of $t_0 \sim 0.63 - 0.89$ s (see Fig. 5.9). Due to the higher acceleration rate of $^3$He (about factor of 2–3 above $^4$He), the fast tail of the energy distribution function populates energies above 1 MeV amu$^{-1}$ faster than for $^4$He. This yields a $^3$He/$^4$He $\gg 1$ for small times $t \ll t_0$, $^3$He/$^4$He $\sim 0.1 - 1$ for $t_0 \sim 0.63 - 0.89$ s and asymptotically approaches the coronal abundance ratio for times $t \gg t_0$. The derived time range of $t_0$ corresponds well to the typical time scale of electron acceleration during the impulsive phases of flares of 0.1–1 s as manifest by the shortest hard X-ray peaks (Kiplinger et al., 1984). Since the EF wave field is a direct consequence of the electron acceleration process, the lifetime of the field is expected to be on the same order of magnitude. When electron acceleration ceases, the $^3$He/$^4$He ratio is frozen and no further ion acceleration occurs.

The time integrated total number of $^3$He nuclei being accelerated above 1.3 MeV amu$^{-1}$ in an impulsive flare is about $10^{31}$ particles (Reames et al., 1994). Assuming the flaring area to be of about $10^{17}$ cm$^2$ with a scale height of $10^9$ cm, the proton density to be $n_H = 5 \times 10^{10}$ cm$^{-3}$ and $^3$He/H = $5 \times 10^{-5}$, there are $2.5 \times 10^{32}$ $^3$He ions available in the flaring volume. This yields a percentage of about 4% of the pre-flaring $^3$He population that has to be accelerated above 1 MeV amu$^{-1}$. Computing the value of the $^3$He density $n_3(t_0)$ at time $t_0$ and comparing it to the coronal abundance yields a percentage of about $\sim 3.5 - 10.5\%$. Thus the number of accelerated ions at the required energies is sufficient to explain the observed particle fluxes. The model therefore also accounts for criterion (iii).

The $^3$He acceleration by EF waves also generates ion energy spectra that qualitatively reproduce the observed behavior. According to Moebius et al. (1982) the ob-

![Figure 5.9: Abundance ratio $^3$He/$^4$He of ions above 1 MeV amu$^{-1}$ vs. acceleration time.](image)
Chapter 5. Acceleration and Enrichment of $^3$He by EF waves

Figure 5.10: Particle energy spectrum in MeV amu$^{-1}$ at time $t_0 = 0.76$ s. $^3$He is harder than $^4$He, qualitatively reproducing the spectrum observed by Moebius et al. (1982).

The observed spectrum of $^3$He is harder than the spectrum of $^4$He in the energy range of $0.4 - 4.0$ MeV amu$^{-1}$ which is a feature of all $^3$He rich periods in their study. The according particle densities in dependence on energy per nucleon resulting from the model presented herein are depicted in Fig. 5.10, reproducing the observations.

A turnover of the $^3$He spectrum at around $\sim 120$ keV amu$^{-1}$ energies can be seen in Fig. 5.10. This reproduces very nicely the observations by Mason et al. (2000) who reported a turnover at around $\sim 100$ keV amu$^{-1}$. It can be interpreted as the result of the whole $^3$He population being accelerated by EF waves.

Therefore it has been established that the EFI is not only a viable acceleration mechanism for ions during solar flares, it also can account for the observed abundance enhancements of $^3$He over $^4$He, accelerates enough $^3$He nuclei, and reproduces the qualitative behavior of the particle energy spectra.

5.5 Protons and Heavier Ions

The energy density profile may suggest by the same arguments used to explain the enhancement of $^3$He above $^4$He, that H should be enhanced above $^3$He, which is not the case. This is not a real problem of the model because of the very different number densities of the according ion species. The energy in the wave field available for proton acceleration is higher than the energy available for $^3$He. Nevertheless, there are much more protons. Thus the available energy is distributed among a larger number of nuclei and the energy gained per nucleon is smaller for the protons. By the same argument,


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$^3$He should be further enhanced in comparison to $^4$He. This effect was neglected herein since the scope of this paper is to show the ability of fractionation by the accelerator itself.

When discussing the possible acceleration of the protons by EF waves, other effects have to be taken into account. Whereas, due to their low densities, $^4$He and $^3$He can be treated as test particles, this view is not valid for protons and nonlinear effects as the erosion of the spectral wave energy around the gyrofrequency of H have to be considered. Such an analysis requires simulations that are beyond the scope of the presented work.

$^3$He rich events usually are also characterized by enhanced Fe/O ratios. Although the observed enhancement of Fe/O is only about a factor of 8 over the coronal ratio, it is still the strongest enrichment of the heavier ions. Other heavier ions as Ne, Mg, Si are enhanced by factors of 2-3 above O, whereas C, O, N are not enhanced. In this section the enrichment of Fe ions is addressed and briefly discussed in view of the model presented in this work.

In a plasma at a temperature of $10^7$ K, the most probable charge state of $^{56}$Fe is 20+ while $^{16}$O is fully ionized. When resonantly cyclotron accelerated, $^{16}$O should exhibit the same acceleration rate as $^4$He while $^{56}$Fe, following the argumentation for enhanced $^3$He acceleration rate, is much slower heated. As can be seen from Fig. 5.4c the energy available for $^{56}$Fe is much smaller than for $^{16}$O. This behavior was confirmed by carrying out the same simulations for $^{56}$Fe and $^{16}$O as for $^3$He and $^4$He. Although the acceleration via EF waves by cyclotron resonance energizes $^{56}$Fe ions up to energies of 1 MeV amu$^{-1}$ on time scales of $\sim$ 10 s, an enhancement of Fe/O by the same mechanism as for $^3$He/$^4$He presented above cannot be reached. It is in general impossible in the EFI model to explain a simultaneous enhancement of $^3$He/$^4$He and Fe/O when the spectral energy density exhibits a slope as depicted in Fig. 5.4c. The same energy is available for $^4$He and $^{16}$O, but $^3$He has more and $^{56}$Fe has less energy available. The heating rate of $^{56}$Fe will therefore always lie below the $^{16}$O and $^4$He rate, whereas $^3$He always lies above it in the present simple model. Other processes may enter the picture. Non-linear effects as the back-reaction of the wave field in response to energy loss to the ions can have a significant influence on the heating rates. However, the analysis of such effects are beyond the scope of this paper and will not be addressed here.

Similar problems arise for the enrichment of Ne, Mg, Si since their gyrofrequencies also lie below the gyrofrequency of $^4$He. The same arguments as for $^{56}$Fe therefore make an enhancement above $^{16}$O impossible in the present scenario.

The observed heavy ion enrichment is independent of the degree of $^3$He enhancement, although heavier ion enrichments are generally associated with $^3$He (Mason et al., 2000). This is an indication that it may not be the same mechanism that accounts for the acceleration of $^3$He and $^{56}$Fe or other heavier ions.

5.6 Conclusion

A model for the enrichment of $^3$He during impulsive solar flares is presented. The acceleration of $^3$He and $^4$He can be understood as consequence of Electron Firehose (EF)
waves, resulting from an unstable anisotropic electron distribution with $T^e_\parallel > T^e_\perp$. The model does not need additional sources of free energy or plasma properties than those resulting directly from the acceleration of the bulk of electrons. The essential result of the investigation is that EF waves, excited by an anisotropic electron distribution function, can accelerate $^3$He and $^4$He via cyclotron resonance to MeV amu$^{-1}$ energies on time scales of $\sim$ 1 s. In support of this conclusion, the following specific results are found:

1. The EF waves accelerate $^3$He ions via cyclotron resonance. The symmetry of EF waves in $k_\parallel$ greatly enhances the efficiency of the acceleration process with respect to the standard cyclotron resonant wave-particle interaction of a unidirectional propagating wave field. The parallel force driving the ion out of resonance, and therefore limiting the acceleration time in the unidirectional case, cancels out in the time average for the counterpropagating wave field. The total resonant interaction time therefore is strongly increased and the particle can reach high energies.

2. The linear growth of the EF waves self-consistently generates a spectral energy distribution, causing enhanced acceleration of $^3$He with respect to $^4$He. It is found from test particle simulations that the heating rate of $^3$He exceeds the values for $^4$He by about a factor of 2–3.

3. It is shown that this small enhancement in the heating rate of $^3$He above $^4$He already can account for the observed enhancement in $^3$He/$^4$He from the coronal value of $\sim 5 \times 10^{-4}$ up to 0.1 – 1 during impulsive solar flares. Due to the larger heating rates, the high energy tail of $^3$He populates energies above 1 MeV amu$^{-1}$ faster. It reaches the required abundance ratio of $^3$He/$^4$He for an acceleration time scale of $\sim 0.76$ s.

4. The fraction of $^3$He accelerated by the EF waves is large enough to account for the observed particle fluxes. Between 4 – 11% of the coronal $^3$He population in the flare is accelerated above 1 MeV amu$^{-1}$.

5. The acceleration model reproduces the qualitative behavior of the $^3$He energy spectrum with respect to the $^4$He energy spectrum. As observed by Moebius et al. (1982) the $^3$He spectrum is generally harder than the $^4$He spectrum. Moreover, a turnover in the $^3$He energy spectrum around $\sim 100$ keV amu$^{-1}$ is obtained, reproducing the observations by Mason et al. (2000).

EF waves are an inherent property of stochastic electron acceleration by transit-time damping of cascading fast-mode waves. Thus, their ability of accelerating $^3$He and $^4$He supports the scenario of stochastic acceleration of the bulk electrons in impulsive flares.
Chapter 6

Summary and Outlook

6.1 Summary

Several elements of the particle acceleration process in solar flares have been investigated during this thesis. The analysis contains different approaches including observations, analytic calculations and computer simulations.

Joint observations by the Phoenix-2 spectrometer in Zurich and the Radioheliograph in Nançay (NRH) were used to analyze the spatial geometry of metric type III radio emission associated with narrow band metric spikes. Spectral information yielded the temporal association of the events while the radioheliograph provided spatially resolved data at different frequencies. The findings of this investigation support the view of metric narrow band radio spikes either to be a signature of the acceleration process itself or being emitted in close proximity to the accelerator. The observed location of the accelerator is high in the corona, consistent with the generic flare scenario with acceleration taking place above loop-like magnetic structures. The cause of the spike emission still needs to be investigated and its potential use in exploring particle acceleration to be assessed.

The theoretical part of the thesis contains analysis in linearized kinetic theory, analytic modeling of wave-particle systems and numerical test particle simulations. Various aspects of an anisotropic electron distribution function with $T_{\parallel} > T_{\perp}$ were addressed. Such an electron distribution is produced by most of the possible acceleration mechanisms and is therefore a stable feature of the flare electron acceleration process. It has been established that such an anisotropy leads to the Electron Firehose instability (EFI) under solar flaring conditions: the electrons become unstable to a left-hand circularly polarized electromagnetic wave mode, propagating at around the proton cyclotron
The response of the plasma components to this wave mode has been investigated in two steps: i) the response of the electron population to the instability and its evolution in velocity space and ii) the behavior of $^3$He and other heavier ions in the presence of an EF wave spectrum.

The first step in the analysis was focused on the pitch-angle development of the electrons. Drift kinetic theory has been applied in an analytic approach and its range of validity was established by comparison with numerical test particle simulations. The results indicate that the EFI is capable of pitch-angle scattering the electrons to higher perpendicular velocities, delivering a crucial ingredient to many electron acceleration processes as an intrinsic feature.

The problem of $^3$He enrichment during impulsive flares has been addressed by a full analytic treatment and test particle simulations. A new model for $^3$He acceleration and enhancement above $^4$He was found. EF waves excited by the EFI accelerate $^3$He nuclei to the observed energies of $\geq 1$ MeV amu$^{-1}$. Moreover, the spectral properties of the EF waves yield to an enhanced acceleration of $^3$He above $^4$He, reproducing abundance enhancements that are observed in the solar wind.

Not only is the EFI a result of flare electron acceleration, it also provides two major ingredients to the acceleration process: on the one hand it isotropizes the electron population and therefore enables efficient acceleration and on the other hand, the excited EF waves accelerate and enhance $^3$He to the observed values.

### 6.2 Outlook

The observations presented in Chapter 2 are an important way to interpret dynamic spectra by context imaging. With increasing temporal and spectral resolution, new spectrometers rather complicate the interpretation of dynamic spectra without providing spatially resolved data. An upcoming ground based Frequency Agile Solar Radiotelescope (FASR) (Bastian et al., 1998) will open the opportunity to observe radio emission spectrally and spatially resolved, allowing us to do the analysis of Chapter 2 with one single instrument at many more frequencies. In addition, new space born observatories, such as RHESSI (Lin & Dennis, 2002) and the upcoming Solar STEREO mission, eventually will allow first-hand witnesses of the particle acceleration process on the Sun by simultaneous imaging, i.e. in 3D, and spectrally resolved observations in the UV and X- to gamma-ray range.

The investigation of instabilities of anisotropic plasmas with $T_\perp < T_\parallel$ is an intriguing topic and needs to be expanded in several directions. First of all, and probably the most important one, is a close investigation of the obliquely propagating EF mode that is described in Chapter 3. A full particle simulation is necessary to investigate the competitive behavior with the parallel and quasi-parallel EF mode. Since it is a zero frequency mode it is also unclear, what the response of the electron population to the excitation of this mode is. An analysis similar to the one presented in Chapter 4 is a first step, but in both cases a full particle simulation including wave particle interaction
will be essential. The advantage of the presented test particle simulation lies in the reduced amount of data and the increased traceability of the particles and the relevant physical processes whereas the full treatment allows to see the plasma as a composite of wave and particles. Both approaches are essential for a full understanding.

The counteracting processes of energy input by transit-time damping of fast-mode waves and energy loss due to the EFI of the electron population needs investigation. The time scales of the build-up of the anisotropy by transit-time damping and the erosion and pitch-angle scattering have to be compared by means of computer simulation. The question whether the resulting instability remains at a marginal state or is driven very strongly is yet unanswered and is a topic to future research.

The presented model for the acceleration and enrichment of $^3\text{He}$ has several aspects that need further investigation. As in the above, the influence of the particles on the wave field and vice versa has to be investigated in full detail. While the test particle simulations are a good approximation for the minority ions (as $^3\text{He}$ or Fe), the influence of the protons and the resulting erosion of the wave spectrum needs to be investigated in more detail. A simultaneous coexistence of interacting electrons and ions needs to be simulated in order to see the competitive behavior of erosion and build-up of the EF wave spectrum.

An improvement of the test particle code developed for the analysis herein is a future step. Test particle simulations are predestined for implementation on parallel computing facilities: each particle can be treated independently and therefore be advanced on a single node. The version of the code used here did not distribute the particles on the computing nodes. Ensembles of particles were allocated by hand and the usage of the Asgard computing facility at the ETH Zurich was therefore not optimized. A future version of the code could simulate the evolution of the whole population of electrons and not only sample a few regions in velocity space.
Appendix
Appendix A

IDLWhamp: A GUI to WHAMP

A.1 Introduction

The WHAMP computer program solves the dispersion relation for waves in a magnetized plasma. The original WHAMP code was developed and written by Kjell Rönnmark (1982), but the program that is used by IDLWhamp is a modified version by Rezeau & Bélmont (1998). In the following, WHAMP always refers to the new version of the program. Although the program has been greatly improved and debugged, it still provides the user only with a command line interface for data input and output. Hence, while working with WHAMP the user always needs additional software for extracting data from raw output files and for visualization of the obtained data. In order to facilitate the working with the WHAMP code it is reasonable to combine the additional software and the WHAMP core program in a single software package. There were several attempts towards this goal, such as ‘xwaves’ (Littlefield, 1993) or ‘xwhamp’ (Scales et al., 1991), x-windows wrappers for WHAMP, available on the Internet. Both did not meet the needs of the author while working with WHAMP. It therefore was decided to program another package called IDLWhamp. The GUI is based on IDL for it is widely used in the plasma physics and astrophysics community and delivers a suitable programming environment for GUIs. The user still can use the binary of WHAMP in the command line mode in the same version as distributed by Rezeau & Bélmont (1998) besides a slight change in the data output file.

The version of IDLWhamp described herein is Version 2.2 released in November 2001. A short description of the physics of WHAMP follows in the next section. The GUI is described in Sect. A.3. Section A.4 describes the working with IDLWhamp and gives some hints for finding solutions and a short summary concludes this work.

A.2 Physics of WHAMP

The WHAMP code solves for the real and imaginary angular frequencies of the wave mode (the solution of the plasma dispersion relation) as a function of $k_\perp$ and $k_\parallel$, the components of the wave vector perpendicular and parallel to a background magnetic
Appendix A. IDLWhamp: A GUI to WHAMP

There may be multiple solutions (branches) of the dispersion relation at a given value of $k_\perp$ and $k_\parallel$. In that case, the particular solution that WHAMP finds will be determined by the initial guess for the real part of the angular frequency, which causes the code to lock onto a particular solution (branch).

The plasma may be composed of a variable number of up to six plasma species. The total distribution function will be the sum of the distribution functions for all the individual species. The distribution function for each species is considered to be uniform (or homogeneous) in space, but to vary not only with the magnitude of the velocity but also with its direction (see Eq. A.4). Thus each distribution function is a function of $v_\perp$ and $v_\parallel$, the components of the particle velocity perpendicular and parallel to the background magnetic field.

Each species is described by its density, temperature, particle mass, anisotropy and drift velocity along the magnetic field. The following is a short summary of the original report by Rönnmark (1982).

The plasma dispersion function can be written as

$$\text{det} \left( (\mathbf{k}^2 - \mathbf{k}_0^2) \frac{\omega^2}{\omega_0^2} - \mathbf{\tilde{\varepsilon}}(\omega, \mathbf{k}) \right) = 0, \quad (A.1)$$

where, according to linearized kinetic theory, the dielectric tensor $\mathbf{\tilde{\varepsilon}}(\omega, \mathbf{k})$ is given by

$$\mathbf{\tilde{\varepsilon}}(\omega, \mathbf{k}) = 1 - \frac{\omega_p^2}{\omega^2} \left\{ 1 - \sum_j \sum_{n=-\infty}^{\infty} \int \mathbf{\Pi} \sum_j \frac{n\Omega_j}{v_{\perp j}} \frac{\partial}{\partial v_{\perp j}} + \frac{k_\parallel}{v_\parallel} \frac{\partial}{\partial v_\parallel} f_j^0 \right\}. \quad (A.2)$$

The gyrofrequency of the $j$th species is given by $\Omega_j = q_j B/(cm_j)$ and $\omega_p$ denotes the plasma frequency defined as $\omega_p = \sqrt{\sum_j \omega_{p,j}^2}$. The tensor $\mathbf{\Pi}$ is given by the matrix

$$\mathbf{\Pi} = \left( \begin{array}{ccc} \left( \frac{n\Omega_j}{k_\perp} v_{\perp j} J_n \right)^2 & -i \frac{n\Omega_j}{k_\perp} v_{\perp j} J_n' & \frac{n\Omega_j}{k_\perp} v_\parallel J_n^2 \\ -i \frac{n\Omega_j}{k_\perp} v_{\perp j} J_n' & \left( v_{\perp j} J_n' \right)^2 & -i v_{\perp j} \frac{v_\parallel}{v_{\perp j}} J_n' \\ \frac{n\Omega_j}{k_\perp} v_\parallel J_n^2 & i v_\parallel \frac{v_\parallel}{v_{\perp j}} J_n' & \left( v_\parallel J_n' \right)^2 \end{array} \right). \quad (A.3)$$

where the argument of the Bessel function $J_n$ is $k_\perp v_{\perp j}/\Omega_j$. $f_j^0(v_\parallel, v_{\perp})$ in equation (A.2) denotes the zero order distribution function in velocity space of the particle species $j$. The most general form a particle distribution function in the WHAMP code can have is given by

$$f_j^0(v_\parallel, v_{\perp}) = \frac{1}{(\sqrt{\pi} v_{\parallel j}^3)} \exp \left( -\left( \frac{v_\parallel}{v_{\parallel j}} - v_{\parallel j}^2 \right)^2 \right) \times \left[ \frac{\Delta_j}{\alpha_{1j}} \exp \left( -\frac{v_\parallel^2}{\alpha_{1j} (v_{\parallel j}^2)} \right) + \frac{1 - \Delta_j}{\alpha_{1j} - \alpha_{2j}} \times \exp \left( -\frac{v_\parallel^2}{\alpha_{1j} (v_{\parallel j}^2)} \right) - \exp \left( -\frac{v_\parallel^2}{\alpha_{2j} (v_{\parallel j}^2)} \right) \right] \quad (A.4)$$
Appendix A. IDLWhamp: A GUI to WHAMP

This is the original notation used in Rönnmark (1982). The \( \alpha_{ij} \) parameter is the temperature anisotropy \( \alpha_{ij} = T_{\perp i} / T_{\parallel j} \) of the \( j \)-th distribution function. \( \Delta_j \) and \( \alpha_{2j} \) define the depth and size of a possible loss-cone.

The integral over velocity space in Eq. (A.2) is evaluated by the following two relations

\[
\Lambda(\lambda_j) = \int_0^\infty J_n \left( \frac{k_{\perp} v_{\perp}}{\Omega_{\lambda_j}} \right) \exp \left( -\frac{v_{\perp}^2}{(v_{\text{th}}^j)^2} \right) \frac{2v_{\perp}}{(v_{\text{th}}^j)^2} \, dv_{\perp},
\]

where \( \lambda_j = 1/2(k_{\perp} v_{\text{th}}^j / \Omega_{\lambda_j})^2 \), and

\[
Z \left( \frac{\omega - n \Omega_j}{k_{||} v_{\text{th}}} \right) = \frac{1}{\sqrt{\pi}} \int_{-\infty}^{\infty} \exp \left( -\frac{v_{\parallel}^2}{(v_{\text{th}}^j)^2} \right) \frac{1}{k_{||}} \, dv_{||},
\]

where \( \Lambda_n(\lambda_j) = e^{-\lambda_j} I_n(\lambda_j) \), \( I_n \) is a modified Bessel function of order \( n \), and \( Z \) is the plasma dispersion function (Fried & Conte, 1961).

By introducing a Padé approximant for the plasma dispersion function the infinite sums of modified Bessel functions which appear in Eq. (A.2) may be reduced to a summable form. The resulting expression for \( \bar{\epsilon}(\omega, \vec{k}) \) is valid for all real \( \vec{k} \) and suited for numerical evaluation.

Additional information about the general theory underlying the program and the FORTRAN code comprising the WHAMP program are described in Rönnmark (1982).

A.3 The GUI

When starting up IDLWhamp, IDL forms a new window (see Fig. A.1) in the upper left corner of the screen. The window contains a menu bar on top (Figure A.1, A) through which all features of IDLWhamp are accessible, the most often used actions in form of a tool bar (Figure A.1, C), the first page of WHAMP inputs (Figure A.1, D) and a dialog field (Figure A.1, E) keeping the user informed of the current status of IDLWhamp. The inputs on the main window contain the parameters for the solving iteration: the range in \( |k| \), increment of \( |k| \), the range in angle \( \Theta \), the angle between \( \vec{k} \) and the direction of the background magnetic field, increment in \( \Theta \) and the initial guess for the complex angular frequency. For information on how to choose starting values for the iteration see Sect. A.4.

The top line of the IDLWhamp main window (Figure A.1, B) always displays the path and name of the current dataset. The pull down menu is divided into seven categories named 'File', 'Edit', 'Navigate', 'Plot', '3DView', 'Data' and 'Help'. The 'File' and 'Edit' menus contain, besides the usual functions, some special items that are described below in the tool bar section. The 'Navigate' menu is a dynamic list of the ten last datasets loaded during a session. It allows the user to switch between dataset in an emacs-like style. The '3DView' menu contains buttons for the 3D plotting routines.
Appendix A. IDLWhamp: A GUI to WHAMP

Figure A.1: Image of the IDLWhamp main window. All the iteration parameters are changed via this window.

displaying the distribution functions of single plasma species. In order to use them, only one plasma species needs to be activated in the 'Enter Setup' window. The 'Data' menu contains the link to a plasma property calculator allowing quick information on the properties (plasma beta, gyrofrequency, etc.) of a chosen species and another button for viewing the raw ASCII WHAMP data file. The functionality of the rest of the pull down menus is self describing and needs no explanation.

The tool bar mirrors several functions of the menu bar. Some of them are special to IDLWhamp and shall be explained here in detail. 'Lock Setup' is a function that prevents the user from accidentally altering a dataset. If a setup is locked, no inputs of any kind are possible anymore, only plotting routines and unlocking are allowed. 'Import Setup' allows the user to import an already existing setup into a newly created setup. 'Print Setup' prints out the whole setup (name, iteration parameters, plasma parameters). The rest of the tool bar items concerns setup handling and plotting and is straightforward.

Instead of typing in the iteration parameters IDLWhamp provides a graphical input interface (Figure A.2) which allows the user to determine the starting values via mouse. A region in the \( \omega, |\vec{k}| \) plane is displayed (Figure A.2, C) and can be shifted to arbitrary values (Figure A.2, D). By choosing one of the plasma species by pressing one of the buttons (Figure A.2, A) the characteristic values for plasma frequency \( \omega \), gyrofrequency \( \Omega \) and several wave vectors (\( k_I \) inertial length, \( k_D \) Debye length, \( k_L \) Larmor radius)
are displayed. Starting values now can be chosen by left clicking on some point in the plane. The range in $|k|$ is determined by a second left click which also starts the analysis. The number of points that should be computed is chosen via the check boxes in the window (Figure A.2, B). Three dialog fields keep the user informed on the current status (Figure A.2, E).

After launching the main window IDLWhamp loads a dataset called 'defaultset'. It is hardcoded in the source of IDLWhamp and therefore cannot be changed permanently. In case of any changes, the next time IDLWhamp is started all values are reset. For the input of the plasma parameters another window pops up, the so-called 'Enter Setup' window. Here the plasma distribution functions are defined by the input of name, charge, density, mass, parallel temperature, perpendicular temperature, drift velocity, loss cone depth and size for each of the plasma species. The setup can be saved manually by choosing the option from the 'File' menu. This creates a file in the setup/ subdirectory of the IDLWhamp tree containing all defining parameters of the current dataset. This file is only used by the GUI.

The 'Enter Setup' windows also displays the resulting charge when the 'Update charge' button is pressed or the window is closed. In case of the resulting charge not being equal to zero a window pops up notifying the user of a problem with charge neutrality. This does not necessarily mean, that charge neutrality is violated. Due to numerical accuracy problems in IDL it sometimes is not possible to correctly compute the resulting charge. In this case it is recommended to check the plasma parameters 'Density' and 'Charge'.

By choosing the 'Run Analysis' action, IDLWhamp first saves the current setup and

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**Figure A.2:** Image of the IDLWhamp main graphical input interface.
creates a parameter file in the subdirectory par/ which is used as an input file for the
WHAMP code. The WHAMP code is started by IDLWhamp spawning a call to the
executable. The output of the WHAMP code is copied to a file in the data/ subdirectory.

Three files are therefore needed for a full dataset. After having run an analysis,
the data can directly be plotted by choosing one of the options from the 'Plot' menu.
The options are grouped in four different subsections and a general plot window: 'Wave
Dispersion', 'Electric Polarization', 'Magnetic Polarization', 'Resonance Factors' and
'General 2D Plots'. The last one offers the opportunity to choose from a list containing
all possible output data from WHAMP for x-axis and y-axis. The former ones offer
nicely formatted plots for the most important and frequently used data and derived
quantities. There are slight differences in the plot windows but in the following we only
want to describe the specialized plot window in detail.

![Plot 2D window](image)

**Figure A.3:** Image of the IDLWhamp plot window.

The plot window (see Fig. A.3) offers a wide diversity of interactive data analysis.
Interactively obtained data points can be written to an ASCII file in two column format
by pressing the right mouse button (Figure A.3, D). Information on electro magnetic
field strengths and other wave parameters are obtained by pressing the middle mouse button. The cross hair can be fixed in x-direction by pressing the left mouse button in order to make an accurate readout of data points easier. In a field at the bottom left of the plot window the current position of the cross hair is continuous displayed in data coordinates (Figure A.3, B). At the bottom of the 'Plot' window there is a menu of check boxes that allows the choice of a plasma species for the 'Resonance Factor' plots (Figure A.3, C).

During the solving process the accuracy of the iteration can drop below a certain tolerance value. IDLWhamp therefore sometimes underlays the plot with blue and/or red color in uncertain regions (dark gray shade in Fig. A.3). A detailed description of the color coding follows in the next section (Note: color coding is not present in the 'General 2D Plots' window!).

A.4 Hints for using IDLWhamp

If an analysis is run with IDLWhamp (by choosing 'Run Analysis' from the pull down menu 'File' or by pressing the 'GO!' button in the tool bar of the main window) it might happen, that no data is displayed in the plot window when opened and the message 'Analysis running...' was not followed by 'Analysis finished' in the dialog field. In this case the WHAMP code did not produce any output (the WHAMP output file can be checked by choosing 'View Raw Data' from the 'Data' menu). This means that some of the input parameters are not accepted by the WHAMP code: charge neutrality etc. have to be checked to be correct. Most possible wrong inputs should be intercepted by IDLWhamp but there still can be some gaps.

When viewing data plots it sometimes happens that regions of the plotting area are under laid with blue or red color. The color coding mirrors two kinds of 'error' messages from WHAMP. Blue stands for the message 'too damped' and red stands for 'no convergence'.

The latter means that WHAMP could not lock-on to a solution of the dispersion relation. The process of finding a solution involves an iterative calculation, and the process is producing a diverging result. This could mean that the initial frequency and/or the starting value for $|k|$ is bad. Dispersion codes tend to be rather sensitive to good initial guesses in order to converge. Knowing a good frequency guess will require either a previous successful run (like the run which results from the default parameters) from which you can gradually change the parameters, or some knowledge of the physics (i.e. results obtained from cold plasma physics or warm plasma solutions). Also the units used should be checked. Information on the requested units for each plasma parameter can be found by pressing the according button in the 'Enter Setup' window (i.e. by pressing the button 'Drift Vel.' a small window pops up saying 'The drift velocity in units of the actual thermal velocity', see also the Appendix).

Blue color indicates that the damping rate (imaginary part of the angular frequency) is too large (WHAMP recognizes that in such a case the results are not good – the results may still be correct, but they cannot be trusted anymore) In that case, WHAMP did
lock-on a solution of the dispersion relation, but the results are not good. Different starting values for $|k|$ and/or $\Theta$ should therefore be used. It might be possible to lock-on to a region of the dispersion surface which leads to good results.

If only parts of the plotting region are under laid with blue (red), WHAMP successfully has found the angular frequency for several combinations of $|k|$ and $\Theta$. But sometimes, the code will be unable to solve for certain combinations of $|k|$ and $\Theta$. As described above the two reasons why this can occur are either that the damping rate was too large (sometimes because the angular frequency is approaching a gyrofrequency), or that the code was simply unable to converge on a result (sometimes because the angular frequency is increasing too steeply – changing the $|k|$ and/or $\Theta$ increment can sometimes help with this problem). When WHAMP does not give a result for a value of $|k|$ and $\Theta$, it continues to run, but it skips to the next value of whichever is varied, resulting in red regions in the plot. Transitions over blue(red) regions should not occur in data that should be trusted because it never is ensured that the solver stays on the same branch of solution. Nevertheless it is alright for quick viewing dispersion relations and the search for good initial guesses.

### A.5 Summary

WHAMP is a well known linear dispersion code which has been publicly available for about twenty years now. It has been debugged by many users and the results are very reliable. It therefore is, also due to its public availability, a very helpful tool for all communities involving plasma physics. An extension to the WHAMP code is presented here that does not concern the mathematics and physics of WHAMP, but it provides the potential user with an easy to use and straightforward graphical user interface called IDLWhamp. The greatly improved speed of input and output combined with enhanced computer power of modern platforms allow to explore vast regions of $(\omega, k)$ space in a new way. It has been used also as an efficient tool for planning and interpretation of numerical simulation at the particle level (Messmer, 2002). It also allows a concise handling of datasets and includes plotting routines for the most important results of WHAMP. Every possible output parameter of WHAMP can easily be accessed, plotted and analyzed directly from within the program. Obtaining solutions with WHAMP therefore became easier and much faster.
## A.6 Units

The following table displays the units used in IDLWhamp. They differ from the original normalizations used in WHAMP, described in Rönnmark (1982).

<table>
<thead>
<tr>
<th>Name</th>
<th>Symbol</th>
<th>Units</th>
</tr>
</thead>
<tbody>
<tr>
<td>Modulus of wave vector</td>
<td>$</td>
<td>\mathbf{k}</td>
</tr>
<tr>
<td>Angle $\mathbf{k}, B_0$</td>
<td>$\Theta$</td>
<td>deg</td>
</tr>
<tr>
<td>Frequency, Growth rate</td>
<td>$\omega, \gamma$</td>
<td>rad/s</td>
</tr>
<tr>
<td>Background magnetic field</td>
<td>$</td>
<td>\mathbf{B}_0</td>
</tr>
</tbody>
</table>

### Plasma Parameters of Species $j$

<table>
<thead>
<tr>
<th>Name</th>
<th>Symbol</th>
<th>Units</th>
</tr>
</thead>
<tbody>
<tr>
<td>Charge</td>
<td>$q_j$</td>
<td>e</td>
</tr>
<tr>
<td>Density</td>
<td>$n_j$</td>
<td>$cm^{-3}$</td>
</tr>
<tr>
<td>Mass</td>
<td>$m_j$</td>
<td>$m_P$</td>
</tr>
<tr>
<td>Temperatures</td>
<td>$T_{ij}^\parallel, T_{ij}^\perp$</td>
<td>K</td>
</tr>
<tr>
<td>Drift velocity</td>
<td>$v_D^j$</td>
<td></td>
</tr>
<tr>
<td>Loss cone depth, size</td>
<td>$\Delta_j, \alpha_{2j}$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$v_{th}^j$</td>
<td></td>
</tr>
</tbody>
</table>
Curriculum Vitae

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  Gunnar Paesold and Arnold O. Benz,

- Acceleration and enrichment of $^3\text{He}$ in impulsive solar flares by Electron Firehose waves
  Gunnar Paesold, Reinald Kallenbach, Arnold O. Benz,

- Spatial analysis of solar type III events associated with narrowband spikes at metric wavelengths
  Gunnar Paesold, Arnold O. Benz, Karl-Ludwig Klein, Nicole Vilmer,

- Electron Firehose instability and acceleration of electrons in solar flares
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• Acceleration and enrichment of $^3$He in impulsive solar flares by Electron Firehose waves
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• Spatial analysis of solar type III events associated with narrowband spikes at metric wavelengths
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  *Geophysical Research Abstracts, Vol. 2*
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