The electrical asymmetry effect in capacitively coupled radio-frequency discharges

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Abstract
We present an analytical model to describe capacitively coupled radio-frequency (CCRF) discharges and the electrical asymmetry effect (EAE) based on the non-linearity of the boundary sheaths. The model describes various discharge types, e.g. single and multi-frequency as well as geometrically symmetric and asymmetric discharges. It yields simple analytical expressions for important plasma parameters such as the dc self-bias, the uncompensated charge in both sheaths, the discharge current and the power dissipated to electrons. Based on the model results the EAE is understood. This effect allows control of the symmetry of CCRF discharges driven by multiple consecutive harmonics of a fundamental frequency electrically by tuning the individual phase shifts between the driving frequencies. This novel class of capacitive radio-frequency (RF) discharges has various advantages: (i) A variable dc self-bias can be generated as a function of the phase shifts between the driving frequencies. In this way, the symmetry of the sheaths in geometrically symmetric discharges can be broken and controlled for the first time. (ii) Almost ideal separate control of ion energy and flux at the electrodes can be realized in contrast to classical dual-frequency discharges driven by two substantially different frequencies. (iii) Non-linear self-excited plasma series resonance oscillations of the RF current can be switched on and off electrically even in geometrically symmetric discharges. Here, the basics of the EAE are introduced and its main applications are discussed based on experimental, simulation, and modeling results.

(Some figures in this article are in colour only in the electronic version)

1. Introduction

Various types of capacitively coupled radio-frequency (CCRF) discharges are frequently used for different applications such as semiconductor manufacturing and the production of solar cells [1, 2]. Each type provides unique features useful for particular applications; CCRF discharges can be geometrically symmetric, i.e. the surface areas of the powered electrode, $A_p$, and the grounded electrode, $A_g$, are identical ($A_p = A_g$), or geometrically asymmetric ($A_p \neq A_g$). In geometrically asymmetric discharges a dc self-bias, $\eta$, is generated via the geometric discharge asymmetry [3–6]. This bias leads to high sheath voltages and high ion energies at the smaller electrode. However, $\eta$ cannot be easily adjusted to control the ion energy. Such discharges are frequently used for etching processes, where high ion energies at the substrate surface are required.

Large area discharges, such as those used for the production of solar cells, are typically geometrically symmetric because of the large aspect ratio of the electrode diameter and gap. In these discharges no dc self-bias is generated geometrically. Consequently, the sheaths adjacent to each electrode are necessarily symmetric and the mean sheath voltages are identical. These constraints are a significant problem for plasma-enhanced chemical vapor
deposition (PECVD) processes, for which high process rates (ion fluxes) and low ion energies at the substrate surface are ideal. In classical large area discharges, the discharge symmetry does not allow realization of these conditions in an optimum way, since high ion fluxes require high driving voltages (or powers), while low ion energies at the substrate surface require low driving voltages.

In contrast to single frequency discharges the ion energy and flux at the electrodes can be controlled separately (to some extent) in discharges driven at multiple, substantially different frequencies [7–14]. Typically, two substantially different frequencies, e.g. 2 and 27 MHz, are used. The idea to obtain separate control of these ion properties in such classical dual-frequency (df) discharges is the functional separation of both frequencies due to $f_{HF} \gg f_{LF}$. The high frequency (HF) voltage amplitude is assumed to sustain the plasma and, consequently, to control the charged particle density and the ion flux. The low frequency (LF) voltage amplitude is assumed to control the acceleration of the ions in the sheaths determining the mean ion energy, without affecting the ion flux. Previous works, however, have shown that this functional separation and, therefore, the separate control properties in classical df discharges is strongly limited by the frequency coupling and the effect of secondary electrons [15–23]. Alternatively, hybrid discharges (capacitive–inductive, capacitive–helicon, capacitive–dc) are used to realize separate control of ion energy and flux at the electrode surfaces [24–30].

All these discharges are typically optimized for the application of interest by either trial and error or complicated numerical models or simulations. Here, we present an analytical model to describe different types of CCRF discharges independent of the discharge geometry and applied voltage waveform based on the non-linearity of the boundary sheaths, which allows us to understand the discharge physics and to extract important plasma parameters directly [1, 12, 19, 31–33]. It yields simple analytical expressions for important plasma parameters such as the dc self-bias, the uncompensated charge in both sheaths, the discharge current, and the power dissipated to electrons.

Based on the model results the electrical asymmetry effect (EAE) [34–44] is understood. The EAE solves many problems of modern plasma processing and provides additional options for process control. We propose a novel class of electrically asymmetric CCRF discharges driven by multiple consecutive harmonics of a fundamental frequency with fixed but adjustable phase shifts between the driving harmonics. By experimental, simulation, and model results we demonstrate that in these discharges: (i) A variable dc self-bias is generated as an almost linear function of the phase shifts between the driving frequencies in both geometrically symmetric and asymmetric discharges. In this way the sheath symmetry in geometrically symmetric discharges is broken and controlled. (ii) Almost ideal separate control of ion energy and flux at the electrodes is realized in contrast to classical df discharges. (iii) Non-linear self-excited plasma series resonance (PSR) oscillations of the RF current can be switched on and off electrically even in geometrically symmetric discharges. We show that driving one electrode with two consecutive harmonics, e.g. 13.56 MHz and 27.12 MHz, is technologically the most simple realization of such an electrically asymmetric discharge. Using multiple consecutive harmonics, however, improves the performance of the EAE.

The paper is structured in the following way. In section 2, the analytical model of CCRF discharges is introduced. In section 3, the basic idea of the EAE, the experimental setup of a prototype of an electrically asymmetric discharge with variable geometrical asymmetry, and the particle-in-cell (PIC) simulation used to investigate the EAE are introduced. Moreover, the advantages of the EAE are discussed in detail: first, the generation of a variable dc self-bias and symmetry breaking in geometrically symmetric discharges is discussed; second, it is demonstrated that the EAE allows an almost ideal separate control of ion energy and flux at the electrodes; third, external control of PSR oscillations of the RF current via the EAE is investigated; finally, conclusions are drawn in section 4.

2. Analytical model of capacitively coupled RF discharges

Our model is an equivalent circuit model of a CCRF discharge including its typical external electrical circuit [31–33, 35, 45] (see figure 1). The external circuit consists of an ideal voltage source (generator) in series with a blocking capacitor $C$ (representing the dc self-bias). The discharge itself consists of two non-linear capacitors representing the sheaths adjacent to both electrodes and a series of an inductance (electron inertia) and resistance (electron–neutral collisions) representing the plasma bulk. The bulk part is based on the assumption that the current in the plasma bulk is purely conduction current, since $\omega_{RF} \ll \omega_{pe}$, where $\omega_{RF} = 2\pi f_{k}$ with $f_{k}$ being the

![Figure 1. Equivalent circuit model of a CCRF discharge.](image-url)
the highest of the $k$ driving frequencies and $\phi_{pe}$ the electron plasma frequency. Thus, the following voltage balance describes capacitive RF discharges:

$$\tilde{\phi}_- (t) = \tilde{\phi}_{C} [q(t)] + \tilde{\phi}_{sp} [q(t)] + \tilde{\phi}_q [q(t)] + \tilde{\phi}_b [q(t)].$$  \hspace{1cm} (1)$$

Here $\tilde{\phi}_- (t) = \phi(t)/U_{tot}$ is the applied RF voltage waveform, $\phi(t)$, normalized by $U_{tot}$. $\phi(t)$ is the sum of $k$ consecutive single frequency harmonic waveforms with fundamental frequency $f$ ($T_{RF} = 1/f$), amplitudes $u_n$, and relative phase shifts $\theta_n$ ($n = 1, \ldots, k$ is an integer). $U_{tot}$ is the sum of the applied harmonics’ amplitudes:

$$\tilde{\phi}_- (t) = \frac{1}{U_{tot}} \sum_{n=1}^k U_n \cos (2\pi n f t + \theta_n),$$

where

$$U_{tot} = \sum_{n=1}^k U_n.$$  \hspace{1cm} (3)$$

In equation (1) all voltages are normalized by $U_{tot}$. This normalization is indicated by the bar. $\tilde{\phi}_{C} [q(t)]$ is the voltage drop across the bias capacitor, $\tilde{\phi}_{sp} [q(t)]$, $\tilde{\phi}_q [q(t)]$ is the voltage drop across the sheath adjacent to the powered and grounded electrode, respectively, and $\tilde{\phi}_b [q(t)]$ is the bulk voltage.

The basic idea of the model is to express all voltages as functions of the normalized positive space charge in the sheath adjacent to the powered electrode, $q(t) = Q(t)/Q_0$, where $Q(t)$ is the unnormalized space charge in this sheath as a function of time and $Q_0 = \Lambda_p \varepsilon_0 \bar{n}_{sp} W_{tot}/T_{RF}$ is a normalization constant. Here, $\varepsilon$ is the elementary charge, $\varepsilon_0$ is the dielectric constant, $\bar{n}_{sp}$ is the spatially averaged ion density in the sheath at the powered electrode and $I_{sp}$ is the sheath integral for this sheath [35]:

$$I_{sp} = 2 \int_0^1 p_s(\xi) \xi d\xi$$  \hspace{1cm} (4)$$

with $\xi = x/s_m$ and $p_s(\xi) = n_s(x)/n_{sp}$. Here, $s_m$ is the maximum width of the sheath adjacent to the powered electrode, while $n_{sp}(x)$ is the ion density as a function of the distance perpendicular to the powered electrode, $x$ ($1 \leq I_{sp} \leq 2$ depending on the ion density profile in the sheath; typically $I_{sp} \approx 1$).

If $U_{tot}$ approximately equals the maximum voltage drop across the sheath at the powered electrode, $Q_0$ will correspond to the maximum space charge in this sheath and the total uncompensated charge in the discharge, $Q_1$, which is almost completely located inside the sheath. Therefore, $q(t)$ changes from a minimum value of about 0 to a maximum value of approximately 1 as a function of time within one period of the fundamental applied frequency. Due to the RF floating potentials $Q_1$ is never completely located in one sheath and, thus, the minimum and maximum values of $q(t)$ are not exactly 0 and 1, respectively.

The next step is to express all terms in equation (1) individually and explicitly by $q(t)$. Then, equation (1) can be solved for $q(t)$ and basically all quantities of interest can be calculated immediately, e.g. sheath and bulk voltages, dc self-bias, currents, and the power dissipated to electrons.

We begin with the normalized sheath voltages, $\phi_{sp} [q(t)]$ and $\phi_q [q(t)]$; various experiments and simulations have shown that the charge–voltage characteristic of the sheath is quadratic to a good approximation [39, 47]:

$$\tilde{\phi}_q [q(t)] = -q(t)^2.$$  \hspace{1cm} (5)$$

This quadratic charge–voltage relation reflects the analogy of the sheath to a non-linear capacitor. This non-linearity is physically caused by the time modulated distance between the capacitor plates, i.e. the sheath width. An example for the charge–voltage characteristic in a geometrically asymmetric ($A_p \ll A_g$) single frequency krypton discharge measured by fluorescence dip spectroscopy is shown in the left plot of figure 2 [46, 47].

The positive space charges in the sheaths at both electrodes are linked via the normalized total uncompensated charge in the discharge, $q_t = Q_1/Q_0$. As $Q_1$ is almost completely located inside the sheaths, $q_1(t) = q(t) + q_{sg}(t)$. Here, $q_{sg}(t)$ is the normalized positive space charge in the sheath at the grounded electrode. Charged particles are created by ionization and are lost by recombination or particle losses to the walls. Only the latter process affects the total uncompensated charge, since positive and negative charges are lost separately. Thus, the dynamics of $q_t$ within one RF period is determined by the continuous loss of positive ions to the electrodes due to the Bohm fluxes, which reduces $q_t$, and the instantaneous loss of electrons to the electrodes at the phases of sheath collapse, which increases $q_t$. The dynamics of the total uncompensated charge in the discharge is shown in the right plot of figure 2 for an argon discharge operated at 27.12 MHz, as an example. $q_t$ is found to be modulated by typically max. 10% around its time average value and can, therefore, be assumed to be constant to a good approximation ($q_t \approx 1$) [43]. Finally,

$$q_{sg}(t) = q_t - q(t).$$  \hspace{1cm} (6)$$

The sheath voltages are linked via the symmetry parameter [35]:

$$\varepsilon = \frac{\phi_{sp,max}}{\phi_{sp,max}} = \left( \frac{\Lambda_g}{\Lambda_p} \right)^{2/3} \bar{n}_{sp} \left( \frac{Q_{mg}}{Q_{mp}} \right)^{1/3}.$$  \hspace{1cm} (7)$$

Here $\phi_{sp,max}, \phi_{sp,max}$ is the maximum voltage drop across the sheath at the powered and grounded electrode, respectively, $\bar{n}_{sp}$ is the spatially averaged ion density in the sheath at the grounded electrode, and $Q_{mg}, Q_{mp}$ is the maximum charge in the respective sheath. $\varepsilon$ is unity for completely symmetric discharges and 0 or $\infty$ for strongly asymmetric discharges. With the help of the symmetry parameter the following expression for the voltage drop across the sheath at ground is found [35]:

$$\tilde{\phi}_{sg}(t) = \varepsilon (q_t - q(t))^2.$$  \hspace{1cm} (8)$$

Next, we discuss the normalized voltage drop across the blocking capacitor, $\tilde{\phi}_C$, of capacitance $C$. Generally, for a linear capacitor the following relation holds:

$$\frac{\partial \tilde{\phi}_C}{\partial t} = \frac{I(t)}{CU_{tot}} = -\frac{Q_0}{CU_{tot}} \frac{\partial q(t)}{\partial t}.$$  \hspace{1cm} (9)
Figure 2. Left: charge–voltage characteristic of the sheath adjacent to the powered electrode in a geometrically asymmetric \((A_x \ll A_y)\) single frequency krypton discharge operated at 13.56 MHz at 10 Pa and 8 W measured by fluorescence dip spectroscopy [46, 47]. Right: sheath voltages and total uncompensated surface charge density, \(\sigma_{\text{tot}} = Q_0/A_x\), as a function of time within two RF periods (2\(T_{\text{RF}}\)) in a geometrically symmetric single frequency argon discharge operated at 27.12 MHz, 6.6 Pa, \(d = 2.5\) cm, \(R = 20\%\) and 200 V (PIC simulation).

Here \(I(t) = -Q_0 \frac{\partial q(t)}{\partial t}\) is the conduction current. Integration of equation (9) yields \(\tilde{\phi}_C:\)

\[
\tilde{\phi}_C(t) = -\alpha q(t) - \tilde{\eta},
\]

where \(\alpha = Q_0/(U_{\text{tot}}C) = C_s/C \ll 1\) with \(C_s = Q_0/U_{\text{tot}}\) (typically \(\alpha \approx 10^{-2}\) due to the large capacitance of the blocking capacitor). \(\tilde{\eta}\) is an integration constant, which physically corresponds to the normalized dc self-bias. As \(\alpha \ll 1\) and \(0 \leq q \leq 1\), \(\alpha q \ll q^2\) unless \(q \approx \alpha\). Therefore, the first term in equation (10) can be neglected compared with the sheath voltages. It can also be neglected compared with \(\tilde{\eta}\), if the dc self-bias is higher than a few per cent of \(U_{\text{tot}}\). The error made by neglecting this term is of the order of the floating potentials \((q \ll \alpha\) which are generally small and will be neglected in the frame of this model. Then \(\tilde{\phi}_C\) corresponds to the temporally constant normalized dc self-bias to a good approximation:

\[
\tilde{\phi}_C = -\tilde{\eta}.
\]

Next, we discuss the normalized voltage drop across the plasma bulk; an analytical expression for \(\tilde{\phi}_b[q(t)]\) is found based on the electron momentum balance equation in the plasma bulk [31]:

\[
m_e n_e \frac{\partial \tilde{u}_e}{\partial t} = -e n_e \tilde{E} - m_e n_e v_m \tilde{u}_e.
\]

Here, \(m_e\) is the electron mass, \(n_e\) is the electron density, \(\tilde{u}_e\) is the electron drift velocity, \(v_m\) is the electron–neutral elastic collision frequency and \(\tilde{E}\) is the time-varying electric field. Equation (12) is based on the negligence of \((\tilde{u}_e \nabla)\tilde{u}_e\) and the electron pressure, which is balanced by the static ambipolar field. Furthermore, the collision frequency is assumed to be independent of the electron drift velocity and electromagnetic effects are neglected. For simple geometries (plane parallel electrodes, cylindrical or spherical discharge geometry) the electron velocity in equation (12) can be expressed by a homogeneous electron current, \(I(t) = -e n_e \tilde{u}_e(t) \cdot \tilde{A}\), through a surface \(\tilde{A}\) with \(\partial I(t)/\partial t = -en_e \tilde{A} \cdot \partial \tilde{u}_e/\partial t\) and the following expression for the bulk voltage is derived [31]:

\[
-\nabla \tilde{\phi}_b = \frac{m_e \tilde{A}}{en_e |A|^2 U_{\text{tot}}} \left( \frac{\partial I(t)}{\partial t} + v_m I(t) \right).
\]

Equation (13) can now be integrated over an arbitrary electron path from one sheath edge to the other to yield the bulk voltage \(\tilde{\phi}_b[q(t)]\) [31]. The result is expressed as a function of \(q(t)\) and normalized by \(U_{\text{tot}}\) to derive the following expression for the normalized bulk voltage:

\[
\tilde{\phi}_b[q(t)] = -2\beta^2[\ddot{q}(t) + \kappa \dot{q}(t)].
\]

where

\[
\kappa = \frac{v_m}{\omega_{\text{RF}}},
\]

The dots in equation (14) indicate derivatives with respect to \(\psi = \omega_{\text{RF}} t\) and \(\beta\) is given by the following expression:

\[
\beta = \frac{\omega_{\text{RF}}}{\tilde{\omega}_{\text{pe}}} \left| \frac{A_y}{A_x} \tilde{\phi}_{\text{sp}}\max \right| / A_{\text{surf}} U_{\text{tot}}.
\]

Here \(\tilde{\omega}_{\text{pe}} = \sqrt{e^2 n_e/(\epsilon_0 m_e)}\) is an effective electron plasma frequency. \(\tilde{n}_e\) is an effective bulk electron density defined as \(\tilde{n}_e^{-1} = l^{-1} \int_{s_y}^{s_y} n_e^{-1} \, ds\). \(l = s_g - s_p\) is the bulk length, defined as the difference between the distances of the edges of the grounded and powered electrode sheaths from the powered electrode located at \(s = 0\). \(A = \tilde{n}_e^{-1} l^{-1} \int_{s_y}^{s_y} (n_e |A|)^{-1} \, ds\) is an effective bulk area. Equation (14) corresponds to a voltage drop across an inductance and a resistance.

Finally, the normalized voltage drops across the individual elements of the equivalent circuit representing the discharge and its external electrical circuit, i.e. equations (5), (8), (10), and (14), are collected to express the entire voltage balance (equation (1)) explicitly as a function of \(q(t)\):

\[
\tilde{\phi}_-(t) + \alpha q(t) + \tilde{\eta} = -q(t)^2 + \varepsilon|q(t) - q(t)|^2 - 2\beta^2[\ddot{q}(t) + \kappa \dot{q}(t)].
\]
Equation (17) describes a large variety of CCRF discharges typically used for applications independent of e.g. the number and values of the applied frequencies, pressures and geometry. In the model, there are four external parameters, \(\alpha, \epsilon, \beta\) and \(\kappa\). These input parameters and the applied voltage waveform must be specified in advance. Furthermore, there are two free parameters in equation (17), the normalized dc self-bias, \(\tilde{n}\), and the normalized total uncompensated charge, \(q_t\). These free parameters are determined by the period averaged compensation of electron and ion fluxes at each electrode.

Typically, \(\alpha \approx 10^{-2}\) and \(\beta \approx 10^{-1}\). Therefore, \(\alpha q(t)\) \((\alpha \ll 1)\) and the voltage drop across the plasma bulk \((\beta^2 \ll 1\) and \(\beta^2 \kappa \ll 1)\) can be neglected as a first approximation in equation (17). However, it must be noted that \(\phi_b[q(t)]\) is essential for the self-excitation of PSR oscillations of the RF current, which will be discussed in section 3.6. In highly electronegative discharges, the bulk voltage might not be negligible due to the low electron density in the plasma bulk, which can increase \(\beta\) significantly. Neglecting \(\phi_b[q(t)]\) simplifies equation (17) from a differential equation to a quadratic algebraic equation in \(\tilde{\phi}\):

\[
\tilde{\phi}_-(t) + \tilde{n} = -q(t)^2 + \epsilon[q_t - q(t)]^2. \tag{18}
\]

\(\tilde{n}\) and \(q_t\) are calculated by considering equation (18) at two distinct times within the RF period, namely the times of maximum applied voltage, \(\phi_{\text{max}}\), and minimum applied voltage, \(\phi_{\text{min}}\) [35]. At these two distinct times, one sheath voltage is maximum, the other sheath voltage is minimum and equation (18) yields

\[
\tilde{\phi}_{\text{max}} + \tilde{n} = \phi_{\text{sp}, \text{max}} + \phi_{\text{sp}, \text{max}}, \tag{19}
\]

\[
\tilde{\phi}_{\text{min}} + \tilde{n} = \phi_{\text{sp}, \text{max}} + \phi_{\text{sp}, \text{max}}. \tag{20}
\]

Note that equations (19) and (20) are not normalized. \(\phi_{\text{sp}}\) and \(\phi_{\text{sp}}\) are the RF floating potentials at the powered and grounded electrode, respectively, which correspond to the minimum sheath voltages of the respective sheath.

First, we derive an analytical expression for the dc self-bias, \(\eta = \tilde{n}U_{\text{tot}}\), by eliminating \(\phi_{\text{sp}, \text{max}}\) and \(\phi_{\text{sp}, \text{max}}\) from the sum of equations (19) and (20), by equation (7):

\[
\eta = \frac{\tilde{\phi}_{\text{max}} + \tilde{\phi}_{\text{min}}}{1 + \epsilon} + \frac{\phi_{\text{sp}, \text{max}} - \phi_{\text{sp}, \text{max}}}{1 + \epsilon} = \frac{\tilde{\phi}_{\text{max}} + \epsilon \tilde{\phi}_{\text{min}}}{1 + \epsilon}. \tag{21}
\]

If the floating potentials are small compared with the applied voltage, the second term of equation (21) will be negligible and only the first term must be taken into account. In this case, the dc self-bias can be calculated simply from the extrema of the applied voltage waveform and the symmetry parameter. Usually the floating potentials can be neglected. However, in multi-frequency discharges with driving voltage waveforms that lead to sheaths similar to dc sheaths, the floating potentials can play an important role [44]. According to equation (21) in symmetric discharges (\(\epsilon = 1\)) driven by a voltage waveform with \(\phi_{\text{max}} = -\phi_{\text{min}}\), e.g. sinusoidal waveforms, there is no dc self-bias, i.e. \(\eta = 0\) V.

Second, we calculate \(q_t = Q_t/Q_0\). If the floating potentials are neglected, which is well justified for typical discharge conditions, \(Q_t\) will be located purely in one sheath at the times of maximum and minimum applied voltage, so that \(\phi_{\text{sp}, \text{max}}/U_{\text{tot}} = -\tilde{q}_t^2\) and \(\phi_{\text{sp}, \text{max}}/U_{\text{tot}} = \epsilon \tilde{q}_t^2\). Then, eliminating \(\eta\) from equations (19) and (20) yields

\[
q_t = \sqrt{\frac{\phi_{\text{max}} - \phi_{\text{min}}}{(1 + \epsilon)U_{\text{tot}}}} \tag{22}
\]

Next, we derive an analytical expression for \(q(t)\) by solving equation (18) for \(q(t)\) [42]:

\[
q(t) = \frac{-\epsilon q_t + \sqrt{\epsilon q_t^2 - (1 - \epsilon)(\tilde{n} + \tilde{\phi}_-(t))}}{1 - \epsilon}. \tag{23}
\]

From the simplified version of equation (21) and equations (22) as well as (23), i.e. from the symmetry parameter and the applied voltage waveform, further plasma parameters can be calculated, i.e. the sheath voltages (equations (5) and (8)), the electron conduction current density \(j_e = -Q_0/\epsilon A_q(\partial \phi/\partial t)\), and the power dissipated to electrons, \(P_e = (j_e^2/\sigma)A_q\), \(\sigma = \tilde{n}_1\epsilon^2/(m_e v_m)\) is the bulk conductivity and \(\tilde{n}_1\) is an effective plasma density in the bulk. The results for \(j_e(t)\) and \(P_e(t)\) are

\[
j_e(t) = j_{e,0} \frac{\sqrt{1 + \epsilon \tilde{\phi}(t)}}{\tilde{\phi}_{\text{max}} - \epsilon^2 \tilde{\phi}_{\text{min}} - (1 - \epsilon^2) \tilde{\phi}(t)}, \tag{24}
\]

\[
P_e(t) = P_{e,0} \frac{(1 + \epsilon) \tilde{\phi}_t^2 (t)}{\phi_{\text{max}} - \epsilon^2 \tilde{\phi}_{\text{min}} - (1 - \epsilon^2) \tilde{\phi}(t)}. \tag{25}
\]

Here \(j_{e,0} = -\omega_P \sqrt{\epsilon \tilde{n}_1 \epsilon^2/(2 I_{Dq})}\) and \(P_{e,0} = j_{e,0}^2 A_q/\sigma = \omega_P \epsilon \tilde{n}_1 \epsilon^2 A_q I_{Dq}/(2 \tilde{n}_1 \epsilon^2 I_{Dq})\). The electron density and the ion flux to the surfaces are typically proportional to the period averaged value of equation (25), \(P_e\) [1, 45]. Based on this model, the excitation dynamics in CCRF discharges can also be described correctly [42].

3. The electrical asymmetry effect

3.1. The idea of the EAE

The dc self-bias, \(\eta\), is an important discharge parameter in CCRF discharges, since it affects the temporally averaged mean sheath voltages. By temporally averaging equation (18) over one period of the lowest applied frequency it can be shown [45] that \(\eta\) corresponds to the difference of the absolute values of the period averaged sheath voltages, \(|\phi_{\text{sp}}|\) and \(|\phi_{\text{sp}}|\):

\[
\eta = \langle |\phi_{\text{sp}}| \rangle - \langle |\phi_{\text{sp}}| \rangle. \tag{26}
\]

Therefore, \(\eta\) affects the mean energy of ions impinging on the electrodes. Classically (without using the EAE), a dc bias can only be generated in geometrically asymmetric discharges \(A_p \neq A_q\) or by applying an additional dc potential. The latter technique does not work for dielectric electrodes or wafers. As the discharge geometry cannot practically be changed during
Figure 3. Normalized driving voltage waveforms as a function of time within one period of the fundamental frequency. Left: electrically asymmetric df discharge \( (f + 2f, U_1^{(2)} = U_2^{(2)}) \) for different values of \( \theta_1 \). Right: single frequency and classical df discharge operated at substantially different frequencies \( (f + 14f, U_1^{(2)} = 4U_2^{(2)}) \).

a plasma process and other parameters affecting the amplitude of \( \eta \), e.g. pressure, driving frequency and voltage amplitude, also affect the ion flux, \( \eta \) could not be used as an adjustable parameter to control the ion energy separately from the ion flux at the electrodes.

However, based on the above model (equation (21)) there is an alternative way to generate a dc self-bias electrically instead of geometrically; even if \( A_s = A_g \) and \( \bar{n}_s = \bar{n}_g \), i.e. \( \varepsilon = 1 \), a dc self-bias will be generated, if the absolute values of the amplitudes of the applied RF voltage waveform are not identical, i.e. \( \phi_{\max} \neq |\phi_{\text{rms}}| \). This is the fundamental idea of the EAE [34, 35].

The most simple way to realize such an RF voltage waveform to generate an electrical asymmetry is driving one electrode by the superposition of a fundamental frequency and its second harmonic with identical amplitudes of both harmonics. Such an electrically asymmetric df discharge corresponds to a special case of equation (2) with \( k = 2 \) and \( U_1^{(2)} = U_2^{(2)} \). Then, \( \theta_1 \) is the relative phase shift between the driving frequencies and \( \theta_2 \) is set to 0°. By tuning \( \theta_1 \) the difference between the absolute values of the extrema of the applied voltage waveform and, therefore, (equation (21)), \( \bar{n} \), and the mean ion energy at the electrodes can be controlled (left plot of figure 3).

For the first time, the EAE allows the breaking of the symmetry of the sheaths electrically [34], i.e. the mean sheath voltages can be made different by changing the applied voltage waveform (equation (26)). In large area CCRF discharges the EAE is the only way to break the sheath symmetry, since these discharges are naturally geometrically symmetric. Thus, in these discharges the EAE is the only way to control the mean ion energies at the substrate surface without changing the applied voltage amplitudes, i.e. without changing the plasma density and ion flux (see section 3.5). This application of the EAE is expected to be most relevant for large area PECVD processes in CCRF discharges such as used for solar cell manufacturing.

### 3.2. Prototype of an electrically asymmetric capacitive RF discharge

The setup of a prototype of an electrically asymmetric CCRF discharge used to investigate the EAE in geometrically symmetric as well as asymmetric df discharges is shown in figure 4. A voltage waveform corresponding to equation (2) with \( k = 2 \) and \( f = 13.56 \text{MHz} \) is applied to the bottom electrode. Two synchronized function generators are used to generate the two driving harmonics (13.56 and 27.12 MHz). The LF and HF voltage waveforms are generated, amplified and matched individually. Electrical filters are used to prevent the other harmonic from penetrating into the amplifier. The phase shift \( \theta_1 \) between the two frequencies is adjusted via the LF generator (\( \theta_2 = 0^\circ \)).

The reactor is a modified GEC reference cell [48]. Both stainless steel electrodes are circular with identical surface areas (radius 5 cm). The plasma is confined between the electrodes by either a glass cylinder (geometric symmetry) or a grounded metal mesh (geometric asymmetry). In the case of plasma confinement by the glass cylinder a small capacitive coupling between the glass and the outer grounded chamber walls slightly reduces the otherwise perfect geometric discharge symmetry. This capacitive coupling effectively enlarges the surface area of the grounded electrode [3, 6, 38]. In the case of plasma confinement by the grounded metal mesh, the surface area of the mesh enlarges the surface area of the grounded electrode and the discharge is geometrically asymmetric. The geometrical asymmetry is varied by changing the electrode gap. The discharge is operated in pure argon.

A high voltage probe is used to measure the time resolved voltage drop across the plasma directly in front of the powered bottom electrode. The temporal average of this voltage yields the dc self-bias \( \eta \) [38]. The ion energy distribution functions are measured by a retarding field energy analyzer (RFEA) [49], which is either implemented into the powered or grounded electrode. A SEERS (self-excited electron resonance spectroscopy) sensor [50] implemented into the
grounded electrode is used to measure the current to the grounded electrode time resolved within the RF period.

3.3. Particle-in-cell simulation

The simulation used to investigate the EAE is a one-dimensional (1d3v) bounded plasma PIC simulation complemented with a Monte Carlo treatment of collision processes (PIC/MCC). At the planar, parallel and infinite electrodes, electrons are reflected with a probability $R$ [51]. Secondary electron emission from the electrodes is neglected. The effect of finite secondary electron emission is discussed in detail in [53]. The neutral gas temperature is taken to be $T_g = 400$ K. The discharge is operated in argon with different electrode gaps $d$. The cross sections for electron–neutral and ion–neutral collision processes are taken from [53–55].

In contrast to the experiment, in the simulation the investigations of the EAE are extended to multi-frequency discharges, i.e. a voltage waveform of the form of equation (2) with $k > 2$ is applied to the bottom electrode with $f = 13.56$ MHz.

The dc self-bias, $\eta$, is determined in an iterative way to ensure that the charged particle fluxes to each of the two electrodes, averaged over one period of the fundamental frequency, balance. Details of the PIC simulation can be found elsewhere [36, 56–58].

3.4. Generation of a variable dc self-bias via the EAE

Figure 5 shows the measured normalized dc self-bias, $\tilde{\eta}$, as a function of $\theta_1$ and the amplitude ratio, $U_2^{(2)} / U_1^{(2)}$, in a geometrically symmetric discharge operated at 13.56 MHz + 27.12 MHz in argon at 100 Pa, $d = 1$ cm, $U_1^{(5)} = 50$ V [41]. In the experiment $U_2^{(2)}$ is varied to change the amplitude ratio. Under these conditions $\varepsilon$ is unity for all combinations of $U_2^{(2)} / U_1^{(2)}$ and $\theta_1$. Although $A_p = A_k$ a dc self-bias of max. 25% of $U_{tot}$ is generated electrically. For a given amplitude ratio the bias can be tuned almost linearly from $–25\%$ to $+25\%$ by tuning $\theta_1$ from $0^\circ$ to $90^\circ$. The strongest bias is generated for $U_2^{(2)} / U_1^{(2)} \approx 1/2$. For $U_2^{(2)} / U_1^{(2)} = 0$ and $U_2^{(2)} / U_1^{(2)} \to \infty$ the bias vanishes, since the discharge is operated as a single LF or HF discharge with $\tilde{\phi}_{max} = |\phi_{min}|$. Similar results are found using the analytical model and the PIC simulation [41].

At lower pressures an even stronger dc self-bias can be generated electrically. At high pressures the sheaths are collisional and $\tilde{n}_e = \tilde{n}_g$ independent of $\theta_1$, i.e. $\varepsilon \approx 1$ for all $\theta_1$ in geometrically symmetric discharges (see left plot of figure 6, [42]). However, at low enough pressures the sheaths are collisionless and a finite bias leads to different mean ion densities in both sheaths due to ion flux continuity inside the
Figure 6. PIC simulation: normalized dc self-bias and symmetry parameter as a function of $\theta_1$ in a geometrically symmetric df discharge operated at 13.56 MHz + 27.12 MHz at 100 Pa, $U_1^{(2)} = U_2^{(1)} = 120$ V, $R = 20\%$, $d = 2$ cm (left plot) and at 2.66 Pa, $U_1^{(2)} = U_2^{(1)} = 315$ V, $R = 20\%$, $d = 6.7$ cm (right plot) [42].

Figure 7. Measured normalized dc self-bias as a function of $\theta_1$ in a geometrically asymmetric discharge operated at 13.56 MHz + 27.12 MHz in argon at 1 Pa, $d = 4$ cm, $U_1^{(2)} = U_2^{(1)} = 100$ V.

Sheaths. This causes $\varepsilon$ to deviate from unity at phase shifts of strong dc self-bias and self-amplifies the EAE (see right plot of figure 6, [35,42]). For the same reason higher driving voltage amplitudes cause a stronger self-amplification of the EAE.

The EAE also works in geometrically asymmetric discharges. Figure 7 shows a measurement of $\bar{\eta}$ as a function of $\theta_1$ in a geometrically asymmetric discharge operated at 13.56 MHz + 27.12 MHz in argon at 1 Pa ($U_1^{(2)} = U_2^{(1)} = 100$ V). In this case the plasma was confined by a grounded metal cylinder with an electrode gap of 4 cm ($A_g > A_p$). As a consequence of the negative dc self-bias generated by the geometric discharge asymmetry the curve shown in figure 7 is shifted to more negative values compared with the geometrically symmetric scenario (figure 6). Nevertheless, the EAE still allows tuning of the bias over a similar range.

The performance of the EAE, i.e. the amplitude of the strongest possible, electrically generated dc self-bias and the range of bias control, can be significantly improved by driving one electrode by a series of $k$ (instead of two) consecutive harmonics with individual harmonics’ amplitudes, $U_0^{(k)}$, and relative phase shifts, $\theta_n$, i.e. a voltage waveform described by equation (2) [44]. Such a voltage waveform corresponds to a finite Fourier series with components $U_n^{(k)}$. In order to reduce Gibb’s ringing [59,60] and to maximize the difference between the absolute values of the extrema of $\phi$ the individual harmonics’ amplitudes should be chosen according to the following formula:

$$U_n^{(k)} = U_0 \frac{k - n + 1}{k}. \quad (27)$$

Here, $U_0$ is a constant, which determines the ion flux and the absolute values of the mean ion energies at the electrodes. The optimum amplitude choice in df discharges of $U_2^{(1)}/U_1^{(2)} = 1/2$ is a special case of this more general criterion for the amplitude choice.

The effect of adding more consecutive harmonics with optimized amplitudes to the driving voltage waveform is shown in the left plot of figure 8; by adding more consecutive harmonics the difference between the absolute values of the extrema of the applied voltage waveform is increased. Therefore, according to equation (21) a stronger dc self-bias is generated. This is shown in the right plot of figure 8 as a result of the PIC simulations (black squares and line) using a driving voltage waveform according to equations (2) and (27) for $\theta_n = 0^\circ \forall n$. Adding more consecutive harmonics increases the absolute value of the strongest possible bias substantially; for example, in a triple frequency discharge the absolute value of the strongest possible bias is 33% of the total voltage amplitude, while it is only 25% in the df case. Adding more consecutive harmonics increases the absolute value of the bias further. For $k > 10$ we observe a saturation of the normalized bias at $|\eta| \approx 40\%$.

This stronger bias with maximum $\bar{\eta}_{\text{max}}$ and minimum $\bar{\eta}_{\text{min}} = -\bar{\eta}_{\text{max}}$ can be tuned from its minimum to its maximum by adjusting only the even phase shifts between the driving frequencies [44].

The PIC simulation results are well reproduced by the full analytical model (equation (21), red line and circles in figure 8 taking the floating potentials as input parameters for the model.
from the simulation. If the floating potentials are neglected, the model does not reproduce the simulation results well for $k > 2$. This is primarily caused by an increase in the difference of the absolute values of the floating potentials as a function of $k$ [44].

The above results for multi-frequency CCRF discharges are obtained from an analytical model and a simulation, in which it is a priori assumed that the applied voltage waveform is perfectly matched to the discharge and all power is deposited into the plasma. Practically, however, matching such voltage waveforms, i.e. minimizing the power reflected by the discharge, is a technological challenge. In principle, this might be achieved by generating and matching each harmonic individually similar to the prototype of an electrically asymmetric df discharge described in section 3.2. However, such a multi-frequency discharge has not yet been realized experimentally.

3.5. Separate control of ion energy and flux via the EAE

The opportunity to control $\eta$ almost linearly by tuning the phase shifts between the driving voltage harmonics can be used to control the energy of ions impinging on the electrodes. Figure 9 shows the ion flux energy distribution functions at the powered and grounded electrode in a geometrically symmetric df discharge operated at 6 Pa with optimized harmonics’ amplitudes as a function of $\theta_1$ (PIC simulation).

The mean ion energy, $\langle E_i \rangle = \int E_i f(E_i) \, dE_i / \Gamma_{ion}$, at both electrodes can be changed by a factor of about 2 in df discharges, while the ion flux, $\Gamma_i = \int f(E_i) \, dE_i$, remains constant within ±10% (figure 10). Here, $E_i$ is the ion energy and $f(E_i)$ is the ion flux energy distribution function. Thus, in contrast to classical df discharges [15–23] almost ideal separate control of ion energy and flux is possible in electrically asymmetric discharges and the role of the two electrodes is reversed electrically by tuning the phase angle between the two driving frequencies from 0° to 90°.

In geometrically asymmetric discharges the control range of $\eta$ is shifted to more negative values due to the geometric discharge asymmetry (figure 7). Therefore, the mean ion energies are generally lower at the larger grounded electrode and higher at the smaller powered electrode compared with the geometrically symmetric case. Nevertheless, separate control of ion energy and flux via the EAE is also still possible.
in geometrically asymmetric discharges. This is verified by measurements of the ion flux energy distribution functions as a function of $\theta_1$ in a geometrically asymmetric discharge (figure 11) under the same conditions as figure 7. Figure 12 shows the mean ion energies and fluxes at both electrodes as a function of $\theta_1$ calculated from the measured distribution functions in this geometrically asymmetric discharge. The opportunity to increase the mean ion energy at the larger grounded chamber wall might provide the basis for a novel technique of wall cleaning by highly energetic ion bombardment instead of chemical cleaning techniques.

If a CCRF discharge is driven by multiple instead of two consecutive harmonics, the maximum electrically generated dc self-bias and the range of bias control are substantially enlarged (figure 8, [44]). This results in a broader control range of the mean ion energy at both electrodes, which can be quantified by the control factor $\chi_i$:

$$\chi_i = \frac{\langle E_i \rangle_{\text{max}}}{\langle E_i \rangle_{\text{min}}}.$$  \hspace{1cm} (28)

Here $\langle E_i \rangle_{\text{max}}$ and $\langle E_i \rangle_{\text{min}}$ are the maximum and minimum mean ion energy at a given electrode, which can be realized by tuning the phase shifts $\theta_n$ between the driving harmonics.

The left plot of figure 13 shows the mean ion energy at both electrodes in a geometrically symmetric discharge driven by three consecutive harmonics (PIC simulation, 10 Pa, $d = 2$ cm, $U_{\text{tot}} = 300$ V). The horizontal lines indicate the maximum and minimum mean ion energy that can be realized in a df discharge driven at 13.56 and 27.12 MHz with $U_1^{(2)} + U_2^{(2)} = 300$ V (optimized amplitudes) and otherwise identical discharge conditions. Obviously, the range of ion energy control is enlarged from a factor of about 2 to a factor of about 3 by adding a third harmonic to the driving voltage waveform. The right plot of figure 13 shows $\chi_i$ as a function of $k$ and demonstrates that the ion energy control factor can be increased to about 9 under these discharge conditions by adding further harmonics.

The question of why the ion flux is constant as a function of the phase shifts between the driving frequencies is answered by the analytical model introduced in section 2 for a geometrically symmetric df discharge operated at identical harmonics' amplitudes and at high pressure ($\varepsilon = 1$) as an example [45]. The left plot of figure 14 shows the power dissipated to electrons, $P_e$, as a function of time within one period of the fundamental frequency and $\theta_1$ calculated by equation (25) under the assumption that $P_{e,0}$ is constant, independent of...
Figure 12. Mean ion energy and ion flux as a function of \( \theta_1 \) measured in a geometrically and electrically asymmetric df discharge (1 Pa, \( U^{(1)}_i = U^{(2)}_i = 100 \text{ V} \)).

Figure 13. PIC simulation results for multi-frequency discharges (argon, 10 Pa, \( U_{\text{act}} = 300 \text{ V} \), \( R = 20\% \), \( d = 2 \text{ cm} \)). Left: \( \langle E_i \rangle \) at both electrodes as a function of \( \theta_1 (\theta_1 = \theta_1 = 0) \) in a geometrically symmetric discharge driven by three consecutive harmonics with optimized amplitudes \([31–33, 39, 45]\). The horizontal lines indicate the max. and min. \( \langle E_i \rangle \) in a df discharge \((13.56 + 27.12 \text{ MHz, optimized amplitudes}) \). Right: ion energy control factor \( \chi_i \) as a function of \( k \) in a geometrically symmetric discharge \([44]\).

\( \theta_1 \). Similar results have been found experimentally and by the PIC simulation. A strong dynamics of \( P_e \) within one period of the fundamental frequency is observed, which changes as a function of \( \theta_1 \). Similarly, the electron impact excitation dynamics change as a function of \( \theta_1 \) \([42]\). These changes are caused by the change in the sheath dynamics determined by the applied voltage waveform. The dynamics of \( P_e \), however, changes in a way that results in a nearly constant power dissipation on time averaging over one period of the fundamental frequency (right plot of figure 14). As the ion flux is determined by the time averaged power dissipated to the electrons, it is almost constant, independent of the phase shifts between the driving frequencies. Similar results are found at low pressures \((\varepsilon \neq 1, [45])\).

The opportunity to separately control the ion mean energy and flux at the electrodes in this way via the EAE might allow the avoidance of standing wave effects in large area CCRF discharges used for deposition processes \([61–64]\). In such discharges, high driving frequencies and low voltage amplitudes are used to generate a high ion flux and low ion energies at the substrate. Due to the high driving frequencies and the large electrode areas, standing wave effects yield significant inhomogeneities across the wafer. The EAE might solve this problem, since it principally allows the use of lower driving frequencies to avoid standing wave effects, while still ensuring a high ion flux and low ion energies by choosing higher driving voltages and electrically adjusting the dc self-bias in a way to ensure low ion bombardment energies at the substrate.

3.6. Control of self-excited plasma series resonance oscillations of the RF current via the EAE

Until now the bulk part of the discharge, i.e. electron inertia and collisions (figure 1), has been neglected. However, electron inertia can lead to the self-excitation of PSR oscillations of the RF current during the sheath collapse due to the sudden change in the sign of the current at this time within the RF period. These oscillations of the RF current typically have frequencies significantly higher than the highest driving frequency, which can enhance ohmic and stochastic electron heating via nonlinear electron resonance heating (NERH) \([31–33, 39, 45]\). If
the bulk part is taken into account, the voltage balance for a CCRF discharge is
\[
\ddot{\phi}(t) + \ddot{\eta} = -q(t)^2 + \varepsilon[q_i - q(t)]^2 - 2\beta^2[q(t) + \kappa \dot{q}(t)].
\]  
(29)

PSR oscillations will be self-excited by the non-linear charge-voltage relation of the sheaths, if (i) collisional damping is low enough, so that the damping of the PSR oscillations is low \((\kappa \beta \ll 1, \text{low pressure})\), (ii) if the discharge is sufficiently asymmetric, i.e. \(\varepsilon \neq 1\) and (iii) if \(\kappa \) is not much greater than \(\beta\). If \(\varepsilon = 1\) (total symmetry) the non-linearity \(q^2\) in equation (29) vanishes and a simple harmonic oscillator equation in \(q\) results. Then \(q\) oscillates like \(\dot{\phi}(t)\) without higher harmonics.

Until now it was believed that the non-linearities of the sheath charge-voltage relation necessarily cancel and PSR oscillations cannot develop in geometrically symmetric discharges, but can only be observed in geometrically asymmetric discharges \((A_x/A_y \neq 1 \rightarrow \varepsilon \neq 1)\). However, the EAE provides the opportunity to also induce a discharge asymmetry, i.e. \(\varepsilon \neq 1\), electrically in geometrically symmetric discharges at low pressures. In this way the self-excitation of PSR oscillations can be controlled by tuning \(\varepsilon\), by adjusting \(\theta_1\) in a geometrically symmetric df discharge driven at \(13.56 + 27.12 \text{ MHz}\). Here, such a discharge operated at \(3 \text{ Pa}, \ U_1^{(2)} = U_2^{(2)} = 1000 \text{ V}\) and \(d = 2.5 \text{ cm}\) in argon is investigated by a PIC simulation as an example. \(\ddot{\eta}, \varepsilon, \) and the electron conduction current in the discharge center as a function of \(\theta_1\) are shown in figure 15. At phase shifts resulting in \(\varepsilon \neq 1\), i.e. \(\theta_1 = 0^\circ\) and \(90^\circ\), high frequency PSR oscillations of the current are observed, while at \(\theta_1 = 51.75^\circ\) (\(\varepsilon \approx 1\)) the PSR oscillations vanish almost completely.

Compared with strongly geometrically asymmetric discharges \((A_x \gg A_y \rightarrow \varepsilon \approx 0)\) the electrically induced discharge asymmetry is comparably small here. Therefore,
the amplitude of the high frequency PSR oscillations of the RF current is relatively small and the enhancement of electron heating by NERI is negligible.

External control of the self-excitation of PSR oscillations of the RF current is also observed experimentally. The left plot of figure 16 shows the RF current to the grounded electrode as a function of time within one period of the fundamental frequency measured by a SEERS sensor in a discharge driven by 13.56 MHz + 27.12 MHz at 3 Pa for different $\theta_1$. Additionally, $\eta$ is measured and $\xi$ is calculated from the measured bias and the amplitudes of the applied voltage waveform using the simplified version of equation (21). Although the discharge is confined by a glass cylinder it is not perfectly symmetric at $\theta_1 = 45^\circ$ due to capacitive coupling between the glass cylinder and the outer grounded chamber walls. Therefore, at a given $\theta_1$, $\xi$ is different in the experiment compared with the simulation. Nevertheless, external control of the PSR oscillations by tuning $\xi$ via $\theta_1$ is clearly observed, i.e. strong PSR oscillations are found, if $\xi$ is significantly different from unity, and weak oscillations are observed, if $\xi$ approaches 1.

These experimental results are reproduced by the analytical model: the right plot of figure 16 shows the numerical solution of equation (29) using $\tilde{\eta}$ and $\tilde{\xi}$ from the experiment (left plot of figure 16) as input parameters for $\theta_1 = 0^\circ$ (black solid line) and $\theta_1 = 90^\circ$ (red dotted line). $q_i$ is calculated from equation (22) using $\xi$ and the extrema of the applied voltage waveform from the experiment. Additionally, the electron density is estimated to be about $10^9$ cm$^{-3}$ under these discharge conditions and $s_m$ is determined from the plasma emission to be about 5 nm, which results in $\beta \approx 0.1$. At 3 Pa the electron–neutral elastic collision frequency in argon is $v_m \approx 1.2 \times 10^8$ s$^{-1}$ resulting in $\kappa \approx 1.4$. These values for $\beta$ and $\kappa$ are used as input parameters for the model here.

4. Conclusions

Based on the voltage balance of a capacitive RF discharge including its typical external electrical circuit and the non-linearity of its boundary sheaths we have presented an analytical model of CCRF discharges, which describes various discharge types, e.g. single and multi-frequency as well as geometrically symmetric and asymmetric discharges. This model yields simple analytical expressions for important plasma parameters such as the dc self-bias, the uncompensated charge in both sheaths, the discharge current and the power dissipated to electrons.

Based on the model’s analytical expression for the dc self-bias, $\eta$, the EAE has been analyzed and explained. The model shows that a dc self-bias can not only be generated via a geometric asymmetry of the discharge chamber but also electrically by driving one electrode with a voltage waveform of which the absolute values of its global extrema are different. Such a voltage waveform can be realized by a series of consecutive harmonics of a fundamental frequency. In contrast to a geometrically generated dc self-bias, $\eta$ can be tuned by adjusting the individual phase shifts between the driving frequencies. In this way the mean energy of ions at the electrodes can be controlled separately from the ion flux in an almost ideal way, the role of both electrodes can be reversed electrically, and the symmetry of both sheaths can be broken and controlled. Therefore, the novel discharge type of electrically asymmetric CCRF discharges is expected to be most relevant for various applications, e.g. large area PECVD processes (solar cell manufacturing).

The most simple realization of an electrically asymmetric discharge is driving one electrode by a superposition of a fundamental frequency and its second harmonic with fixed, but adjustable phase shift, $\theta_1$, between the driving frequencies. Then $\eta$ and the mean ion energies at the electrodes depend almost linearly on $\theta_1$. A prototype of such an electrically asymmetric discharge has been built. Separate control of ion energy and flux in such a discharge via the EAE was verified experimentally and by PIC simulations.

The question why the ion flux is constant as a function of $\theta_1$, was answered by the analytical model; using its result for the power dissipated to electrons it was shown that the dynamics of the power dissipation within one period of the
fundamental frequency changes significantly as a function of $\theta_1$. However, on time averaging the dissipated power and, thus, the ion flux are approximately constant independent of $\theta_1$.

Experimentally, it was demonstrated that the EAE also works in geometrically asymmetric discharges, where separate control of ion energy and flux at the electrodes is also possible. In these discharges $\eta$ is shifted by a constant bias generated by the geometric asymmetry. The mean ion energy at the smaller electrode is generally higher compared with the larger electrode for a given $\theta_1$, but can still be tuned via the EAE.

The absolute value of the electrically generated dc self-bias and the control range of $\eta$ and the mean ion energy at both electrodes can be significantly enlarged. This is achieved by adding further consecutive harmonics of the fundamental frequency to the driving voltage waveform with individual harmonics’ amplitudes according to a criterion defined in section 3.4.

As the symmetry of CCRF discharges can now be controlled electrically, the self-excitation of non-linear plasma series resonance (PSR) oscillations of the RF current can be controlled externally by tuning the phase shifts between the driving frequencies. In an electrically asymmetric discharge we demonstrated experimentally, via PIC simulations, and by the analytical model that the symmetry parameter, $\varepsilon$, can be tuned by adjusting $\theta_1$ at low pressures of a few Pa. In this way PSR oscillations can be switched on ($\varepsilon \neq 1$—asymmetric discharge) and off ($\varepsilon = 1$—total symmetry) via the EAE.

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