Screening of the resistive-wall impedance by a cylindrical electron plasma

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Abstract. The effect of an electron cloud on the longitudinal coupling impedance is studied by idealizing it as a cold and uniformly distributed non-neutral plasma of electrons. The beam pipe is assumed to be of circular cross section with a thick resistive wall and the beam charge is idealized as a uniform disk. The electron contribution to the charge and current densities is obtained from the collective electron response to the beam passage through the pipe. In the presence of the electron background, a general closed formula for the longitudinal coupling impedance is obtained. The screening of the coupling impedance with the density of the electron plasma is studied and discussed for typical parameters in an accelerator beam pipe for the under-dense and the over-dense plasma regions.

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1. Introduction

Collective effects induced by electron clouds (ECs) play an important role in synchrotrons and storage rings for high-intensity proton and heavy-ion beams (see [1] for a recent overview). Electrons can be generated in the beam pipe due to different multipacting processes or due to residual gas ionization. The second mechanism is important for the low- to medium-energy heavy-ion beams in the SIS 18 synchrotron at GSI and in the projected SIS 100. At low energies and for high charge states the beam ions can very effectively ionize the residual gas. In a coasting beam the electrons can then accumulate up to a certain saturation level in the beam potential [2, 3]. Above a certain threshold the accumulated electrons induce two-stream-like instabilities in long bunches. In shorter bunches, the wake field generated by the electrons can induced head–tail type or multi-bunch instabilities. Instabilities caused by clouds of electrons were observed in many synchrotrons and storage rings, e.g. in the KEKB low energy ring (LER) [4], in the CERN proton synchrotron (PS) [5], in the super proton synchrotron (SPS) [6] and in the relativistic heavy ion collider (RHIC) [7]. In the SPS as well as in the LER EC densities of \( n_{ec} = 10^{11} \cdots 10^{12} \text{ m}^{-3} \) have been detected. In [8], it has been proposed to represent the effect of ECs on the circulating ion beam in terms of a transverse impedance. These results were confirmed by simulation studies presented in [9]–[12]. The electron density, however, not only generates a wake-field but it also modifies the impedance coming from coupling of the beam to the resistive wall.

In the present work, we analyze the longitudinal effect of an initially homogenous cylindrical electron plasma of density \( n_{ec} \) filling the interior of the beam pipe representing an idealized model of the EC [6]. This assumption of an undisturbed EC will hold only approximatively. The electrons of the cloud are generated locally, their areas of generation still being under investigation. Whether their lifetime allows for adjusting the electron field self-consistently to the beam, is yet unsettled. Dipole excitations of the EC, as observed, will mostly affect the transverse coupling impedance. Our idealized model, hence, solvable in closed form, will hold as a guidance to the many effects occurring in ECs.

This EC shields the longitudinal beam modes, which, in turn, gives rise to a depletion of the longitudinal coupling impedance. This screening then is calculated as a function of the EC plasma frequency \( \omega_{ec} \), namely, \( \omega_{ec} = \sqrt{n_{ec}e^2/m_ee_0} \), where \( m_e \) is the electron mass, \( e \) is the elementary charge and \( e_0 \) is the permittivity of free space [13, 14]. The electron densities...
predicted for the GSI synchrotrons (see [15]), or measured in different ring accelerators, correspond to electron plasma frequencies of the order of \( f_{cc} = \omega_{cc}/2\pi = 1 \cdots 10 \text{ MHz} \). It turns out that the longitudinal EC effect is governed by the ratio \( \omega_{cc}/\omega \) with \( \omega \) being the excitation frequency.

The paper is organized as follows: in section 2, we present the wave equations for the electromagnetic field components in the presence of a beam perturbation and the uniformly distributed non-neutral plasma of electrons. The radial wave equation for the longitudinal electric field component will be obtained and then solved in section 3. The coupling impedance (resistive-wall plus space-charge) will be obtained in section 4. Numerical examples of typical parameters in an accelerator beam pipe are given in section 4.1 for \( \omega > \omega_{cc} \) and in section 4.2 for the case \( \omega < \omega_{cc} \). Finally, a summary and outlook are presented in section 5.

2. Model equations with EC

Consider the motion of a particle beam in the form of a rotationally symmetric disk of radius \( a \) and total charge \( Q \) in a smooth cylindrical pipe of radius \( b \) and circumference \( L \) (see also [16]). The beam moves with a constant longitudinal velocity \( v_b = \beta c \) along the \( z \)-axis through the cold and uniform EC background of density \( n_{ec} \), which fills the whole pipe up to radius \( b \). The total (effective) current \( \vec{J}(\vec{r}, t) \) and charge \( \rho(\vec{r}, t) \) densities are defined as follows:

\[
\vec{J}(\vec{r}, t) = j_0(\vec{r}, t) + j_e(\vec{r}, t), \quad \rho(\vec{r}, t) = \rho_0(\vec{r}, t) + \rho_e(\vec{r}, t).
\]  

(1)

In following along the lines of Mulser’s oscillator model [17], the current \( j_0 \) and charge \( \rho_0 \) densities in equation (1) are those associated with the streaming motion of the beam particles, whereas \( j_e \) and \( \rho_e \) are the collective current and charge densities due to the coupling between the background electrons and the excited electric fields in the cylindrical pipe. For a cold, collisionless and unmagnetized electron background, we obtain the following simple expression for the effective induced EC current density, namely,

\[
j_e(\vec{r}, \omega) = \frac{i \omega e^2 n_{ec}}{m_e \omega^2} \vec{E}(\vec{r}, \omega),
\]

(2)

\[
\frac{\partial \rho_e}{\partial t} + \vec{\nabla} \cdot j_e(\vec{r}, t) = 0 \quad \rightarrow \quad \rho_e(\vec{r}, \omega) = \frac{e^2 n_{ec}}{m \omega^2} \vec{\nabla} \cdot \vec{E}(\vec{r}, \omega).
\]

Upon using Faraday’s and Ampere’s laws, the wave equations satisfied by the magnetic induction \( \vec{B} \) and the electric field \( \vec{E} \) in non-conducting media with \( \epsilon = \epsilon_0 \) and \( \mu = \mu_0 \) are [16],

\[
\nabla^2 \vec{B}(\vec{r}, t) - \frac{1}{c^2} \frac{\partial^2 \vec{B}(\vec{r}, t)}{\partial t^2} = -\mu_0 \vec{\nabla} \times \vec{J}(\vec{r}, t),
\]

(3)

\[
\nabla^2 \vec{E}(\vec{r}, t) - \frac{1}{c^2} \frac{\partial^2 \vec{E}(\vec{r}, t)}{\partial t^2} = \mu_0 \frac{\partial \vec{J}(\vec{r}, t)}{\partial t} + \frac{\vec{\nabla} \rho(\vec{r}, t)}{\epsilon_0},
\]

excited by the beam are \( E_z(\vec{r}, \omega) \), \( E_r(\vec{r}, \omega) \) and \( B_\theta(\vec{r}, \omega) \). These field components correspond to transverse magnetic (TM) modes. Fourier transforming in time and using \( \omega = k_z \beta c \) results in

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the following total source term:

\[
\vec{J}(\vec{r}, \omega) = \left( \frac{Q}{\pi a^2} e^{ik_z z} \right) \hat{z} + i \frac{e^2 \sigma c}{m \omega} \vec{E}(\vec{r}, \omega),
\]

(4)

\[
\rho(\vec{r}, \omega) = \left( \frac{Q}{\pi a^2 \beta c} e^{ik_z z} \right) \hat{z} + \frac{e^2 \sigma c}{m \omega^2} \vec{\nabla} \cdot \vec{E}(\vec{r}, \omega).
\]

With \( \vec{E}(\vec{r}, \omega) \) having a z-dependence such that \( \vec{E}(\vec{r}, \omega) = \vec{E}(r, \omega) e^{ik_z z} \) as well as \( \rho_0(\vec{r}, \omega) \) and \( j_0(\vec{r}, \omega) \), equation (3) after Fourier transforming in time takes on the following form:

\[
\left( \nabla^2 - k_z^2 \right) E_z(r, z, \omega) + \frac{\omega^2}{c^2} \vec{E}(r, z, \omega) = \frac{i k_z Q}{\epsilon_0 \pi a^2 \beta c \gamma_0^2} \hat{z} - \omega \mu_0 j_z + \frac{\omega^2}{\omega^2} \vec{\nabla} \left[ \vec{\nabla} \cdot \vec{E}(r, z, \omega) \right],
\]

(5)

where \( \nabla^2 \) is the transverse Laplacian operator, and \( \gamma_0 \) is the relativistic factor defined as \( \gamma_0^{-2} = 1 - \beta^2 \). By making use of Maxwell’s curl equations, the relations between the longitudinal electric field component \( E_z(r, z, \omega) \) and the transverse components \( E_r(r, z, \omega) \) and \( B_\theta(r, z, \omega) \) are as follows:

\[
E_r(r, z, \omega) = \frac{-i}{k_z (1 - \beta^2 \epsilon_{ee})} \frac{dE_z(r, z, \omega)}{dr},
\]

(6)

\[
B_\theta(r, z, \omega) = \frac{i \omega \mu_0 \epsilon_{ee}}{k_z^2 (1 - \beta^2 \epsilon_{ee})} \frac{dE_z(r, z, \omega)}{dr},
\]

(7)

where the longitudinal dielectric function

\[
\epsilon_{ee} = 1 - \frac{\omega_{ee}^2}{\omega^2}
\]

(8)

has been introduced. The z-component of \( \vec{\nabla} [\vec{\nabla} \cdot \vec{E}(r, z, \omega)] \) on the right-hand side of equation (5) is rewritten in cylindrical coordinates as

\[
\frac{\partial}{\partial z} \left[ \vec{\nabla} \cdot \vec{E}(r, z, \omega) \right] = \left[ \frac{1}{1 - \beta^2 \epsilon_{ee}} \frac{\partial^2}{\partial r^2} - k_z^2 \right] E_z(r, \omega) e^{ik_z z}.
\]

(9)

Also, using \( \omega = k_z \beta c \), equation (5) for the excited electric field component \( E_z \) becomes,

\[
\left( 1 - \frac{\omega_{ee}^2}{\omega^2} \frac{1}{1 - \beta^2 \epsilon_{ee}} \right) \frac{d^2 E_z(r, \omega)}{dr^2} + \frac{1}{r} \frac{dE_z(r, \omega)}{dr} - \frac{k_z^2 \epsilon_{ee}}{\gamma_0^2} E_z(r, \omega) = \frac{i k_z Q}{\epsilon_0 \pi a^2 \beta c \gamma_0^2} \Theta(a - r),
\]

(10)

where \( \Theta \) stands for the Heaviside unit step function [18]. Within the conducting wall of conductivity \( S \), we have the following equation for \( E_z \),

\[
\left[ \frac{d^2}{dr^2} + \frac{1}{r} \frac{d}{dr} - \frac{k_z^2}{\gamma^2} \right] E_z(r, \omega) = 0, \quad b \leqslant r < \infty,
\]

(11)

where the modified relativistic factor due to the conductivity reads

\[
\gamma = \sqrt{1 - i \mu_0 \omega S \gamma_0^2 / k_z^2}.
\]

(12)
3. Unique solution of the radial wave equations

To solve the homogeneous part of equation (10), we introduce the following parameters:

\[
\alpha = \frac{\omega_{cc}^2}{\omega^2} \frac{1}{1 - \beta^2 \varepsilon_{cc}}, \quad \nu = \frac{\alpha}{2} / \left(1 - \alpha\right),
\]

\[
\sigma_{cc}^2 = \frac{k_z^2 \varepsilon_{cc}}{\gamma_0^2} \frac{1}{1 - \alpha} = \frac{k_z^2 \varepsilon_{cc}}{\gamma_0^2} (2\nu + 1)
\]

and assuming a solution of the form \( E_z(r, \omega) = r^{-\nu} V(r, \omega) \), the homogeneous part of equation (10) reduces to the differential equation of the modified Bessel equation for \( V(r, \omega) \), namely,

\[
\left[ r^2 \frac{d^2}{dr^2} + r \frac{d}{dr} - \left( \sigma_{cc}^2 r^2 + \nu^2 \right) \right] V(r, \omega) = 0.
\]

With \( \sigma = k_z / \gamma \), the overall regular general solutions are expressed in terms of the \( \nu \)th-order modified Bessel functions of first \( (I_\nu) \) and second \( (K_\nu) \) kinds as follows:

\[
E_z(r, z, \omega) = e^{ik_z z} \begin{cases} 
A_1 r^{-\nu} I_\nu(\sigma_{cc} r) - i \frac{Q}{\pi a^2 \varepsilon_0 \varepsilon_{cc} k_z \beta c}, & 0 \leq r \leq a, \\
A_2 r^{-\nu} I_\nu(\sigma_{cc} r) + A_4 r^{-\nu} K_\nu(\sigma_{cc} r), & a \leq r \leq b, \\
A_4 K_0(\sigma r), & b \leq r < \infty.
\end{cases}
\]

Using Maxwell’s curl equations, the electromagnetic field components \( E_i^{(cc, co)} \) and \( B_i^{(cc, co)} \) in the EC plasma and the conducting wall regions are obtained from the corresponding \( E_z^{(cc, co)} \) as follows,

\[
E_i^{(cc)}(r, z, \omega) = -\frac{1}{ik_z (1 - \beta^2 \varepsilon_{cc})} \frac{dE_z^{(cc)}(r, z, \omega)}{dr},
\]

\[
B_i^{(cc)}(r, z, \omega) = \frac{\omega \mu_0 \varepsilon_0 \varepsilon_{cc}}{ik_z^2 (1 - \beta^2 \varepsilon_{cc})} \frac{dE_z^{(cc)}(r, z, \omega)}{dr},
\]

\[
E_i^{(co)}(r, z, \omega) = \frac{\gamma^2}{ik_z} \frac{dE_z^{(co)}(r, z, \omega)}{dr},
\]

\[
B_i^{(co)}(r, z, \omega) = \frac{\omega \mu_0 \varepsilon_0 \gamma^2}{ik_z^2} \frac{dE_z^{(co)}(r, z, \omega)}{dr}.
\]

Accordingly, the radial electromagnetic field component \( E_r \) becomes as follows,

\[
E_r(r, z, \omega) = e^{ik_z z} \begin{cases} 
\frac{\sigma_{cc} r^{-\nu}}{ik_z (1 - \beta^2 \varepsilon_{cc})} A_1 I_{\nu+1}(\sigma_{cc} r), & r \leq a, \\
\frac{\sigma_{cc} r^{-\nu}}{ik_z (1 - \beta^2 \varepsilon_{cc})} \left[ A_2 I_{\nu+1}(\sigma_{cc} r) - A_3 K_{\nu+1}(\sigma_{cc} r) \right], & a \leq r \leq b, \\
i \gamma A_4 K_1(\sigma r), & b \leq r < \infty.
\end{cases}
\]
The integration constants are determined from the boundary conditions at \( r = a \) and \( b \), namely, the continuity of \( E_z \) and \( B_\theta \) at \( r = a \) and \( b \). Upon applying these boundary conditions, we obtain the following integration constants,

\[
A_1 = - \left[ \frac{K_{v+1}(\sigma_{ec}a)}{I_{v+1}(\sigma_{ec}a)} + \frac{K_v(\sigma_{ec}b) - \eta_{ec}K_{v+1}(\sigma_{ec}b)}{I_v(\sigma_{ec}b) + \eta_{ec}I_{v+1}(\sigma_{ec}b)} \right] A_3, \tag{19}
\]

\[
A_3 = \frac{Q a^v \sigma_{ec} a}{i \pi a^2 \epsilon_0 \epsilon_{ec} k_z \beta_c} I_{v+1}(\sigma_{ec}a), \tag{20}
\]

\[
\eta_{ec} = \frac{K_0(\sigma b)}{K_1(\sigma b)} \frac{\omega \epsilon_0 \epsilon_{ec}}{i \gamma (S - i \omega \epsilon_0)} \frac{\sigma_{ec} k_z (1 - \beta_0^2 \epsilon_{ec})}{}, \tag{21}
\]

4. Coupling impedance

The beam longitudinal coupling impedance is obtained as follows \([19]\):

\[
Z_{\parallel}(\omega) = - \frac{1}{Q^2} \int_\Omega E_z^{(v)(\omega)}(r, z, \omega) j_0^*(r, z, \omega) \, d^3r
\]

\[
= - \frac{2L}{Q a^2} \int_0^a \left[ A_1 \frac{I_v(\sigma r)}{r^v} - i \frac{Q}{\pi a^2 \epsilon_0 \epsilon_{ec} k_z \beta_c} \right] r \, dr, \tag{22}
\]

where \( L \) is the circumference of the ring and \( j_0(r, z, \omega) = \frac{Q}{\pi a^2} e^{ik_z z} \) is the beam current density. This integral can be solved to yield the closed form expression

\[
Z_{\parallel}(\omega) = - \frac{2L}{Q a^2} \left[ A_1 \left( \frac{I_{v-1}(\sigma_{ec}a)}{\sigma_{ec} a^{v-1}} - \frac{\sigma_{ec}^{v-2}}{2^{v-1} \Gamma(v)} \right) - i \frac{Q}{2 \pi \epsilon_0 \epsilon_{ec} k_z \beta_c} \right], \tag{23}
\]

\[
Z_{\parallel}(\omega) = i \frac{n Z_0}{2 \beta \gamma_0} \frac{4 \gamma_0^2}{k_z^2 a^2} \left[ 1 - 2^{v} I_1^2(\sigma_{0a}) F \right], \tag{24}
\]

with the constants

\[
F = \frac{K_1(\sigma_{0a}) + K_0(\sigma_{0b}) - \eta[K_0(\sigma b)/K_1(\sigma b)]K_1(\sigma_{0b})}{I_0(\sigma_{0b}) + \eta[K_0(\sigma b)/K_1(\sigma b)]I_1(\sigma_{0b})}, \quad \eta = \frac{\omega \epsilon_0 \gamma_0}{i \gamma (S - i \omega \epsilon_0)}. \tag{25}
\]

For a perfectly conducting wall such that \( \eta \to 0 \), equation (24) reduces into the well-known formula for the space-charge impedance without the EC. The resistive-wall impedance with the EC is defined as the difference between the coupling impedance (space-charge plus resistive-wall) and the impedance of a perfectly conducting wall (space-charge), namely \([16]\),

\[
Z_{\parallel}^{res}(\omega, \omega_{ec}, S) = Z_1(\omega, \omega_{ec}, S) - Z_{\parallel}(\omega, \omega_{ec}, S \to \infty).
\]

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4.1. Numerical examples and discussion

The longitudinal coupling impedance above is calculated for typical parameters in an accelerator beam pipe for different strengths of the EC density. In figure 1 the dependence of the resistive-wall impedance on the EC plasma frequency $\omega_{ec}$ is shown. Both real and imaginary parts of the resistive-wall impedance exhibit a decrease with increasing EC density due to the plasma shielding of the wall, and therefore a reduction in the current induced by the beam in the wall of the beam pipe. Figure 2 shows the same behavior of the real part for different harmonic numbers $n = k_z L / (2\pi)$. The same screening effect at $n = 1$ and 1000 for different $a$ values at fixed pipe radius is shown in figures 3 and 4. Figures 5 and 6 show the imaginary part of the coupling impedance. The total imaginary part consists of a resistive-wall part (very small) and a space-charge $Z_{\parallel}(\omega_{ec} = 0, S = \infty)$ part (dominant). For beam radii $a$ approaching the pipe radius $b$, we observe an enhancement of the space-charge impedance, instead of a reduction. If the beam has the same radius as the pipe, the electrons respond mainly along the longitudinal direction. In this case, the resulting electron wake field leads to anti-screening, as shown in [20].

Finally, figure 7 shows the influence of beam energy on the depletion of the longitudinal resistive wall impedance. Due to the strong dependence on $\beta$ in equation (13) and, hence, on $\sigma_{ec}$, the region where the excitation frequency is smaller than the EC frequency is extended to higher frequencies for high energy machines. For low energy, on the other hand, the EC shields the field of the beam only at very low frequency.
Figure 2. Dependence of the real part of the longitudinal resistive-wall impedance on the ratio of EC plasma frequency to excitation frequency $\omega = n\omega_0$ for different harmonic numbers $n$.

Figure 3. Dependence of the real part of the longitudinal resistive-wall impedance on the ratio $a/b$ for a fixed pipe-radius $b = 10$ cm.
Resistive-wall impedance at $n = 1000$

**Figure 4.** Dependence of the real part of the longitudinal resistive-wall impedance on the ratio $a/b$ for a fixed pipe-radius $b = 10$ cm.

**Figure 5.** Dependence of the space-charge impedance on the ratio $a/b$ for a fixed pipe-radius $b = 10$ cm.
Figure 6. Dependence of the space-charge impedance on the ratio $a/b$ for a fixed pipe-radius $b = 10$ cm.

Figure 7. Fraction of critical EC frequency $f_{\text{crit}}^\text{ec}$ where the longitudinal resistive-wall impedance is half depleted as compared to the excitation frequency for low energy (low $\beta$) and high energy ($\gamma \approx 20$). The dashed lines indicate the measured value of EC density of $10^{12}$ and $10^{11}$ m$^{-3}$ in the SPS, LER and SIS18.
4.2. Surface-wave-sustained modes

Electromagnetic waves cannot propagate in an over-dense plasma if the plasma frequency is larger than the excitation frequency, \( \omega_{ec} > \omega \), i.e. when the beam is completely shielded by the electron plasma \([21, 22]\). Then the waves are reflected at the bulk plasma frequency due to the skin effect and become evanescent waves. Their penetration depth corresponds roughly to the skin depth \( \delta_s \approx c (\omega_{ec}^2 - \omega^2)^{-1/2} \). This may give rise to heating a plasma rather than damping it. The waves then do not travel any more in the radial direction but rather propagate along the plasma surface. The wave energy is then transferred to the plasma by the evanescent wave, which enters the plasma perpendicular to its surface and decays exponentially with the skin depth. Due to the heating process, the real part of the impedance becomes negative. This transfer mechanism allows to support over-dense plasmas with electron plasma frequencies beyond the excitation frequency. For even higher excitation frequencies, \( \omega_{ec}/\omega > \sqrt{2} \), these waves do not propagate any more along the surface but are rather overdamped in the longitudinal direction. Such overdamped surface waves have been studied, e.g. in \([21]\).

Our analysis for a low energy beam (\( \beta = 0.155 \)) shown in figure 8 indicates that one can distinguish three regions of a beam embedded in an over-damped plasma: (i) no waves exist for \( 1 < \omega_{ec}/\omega \leq 1.05 \); (ii) in the intermediate region of evanescent waves, \( 1.05 \leq \omega_{ec}/\omega \leq \sqrt{2} \) the real part of the impedance is negative and the imaginary part is positive, thus capacitive; and (iii) overdamped surface waves with positive real part and negative imaginary part exist for \( \omega_{ec}/\omega > \sqrt{2} \). The latter two regions are separated by a resonance transition. Due to the presence of the beam space-charge the critical value \( \sqrt{2} \) is shifted slightly.
5. Summary and outlook

The modification of the longitudinal coupling impedance by a cold and homogeneous electron plasma has been investigated analytically as a function of the electron plasma density. Real and imaginary parts of the resistive-wall impedance as well as the space-charge impedance exhibit a decrease with increasing EC density due to the plasma shielding of the wall. The impedance shows a peak transition near the quasi-stationary plasma surface wave frequency at \( \omega = \omega_{ec}/\sqrt{2} \). As part of our future work, the present analysis will be extended to include Joule heating of the electrons and the spatial variation of the EC density. Numerical simulations of the longitudinal wake field induced in an electron background will be a next step and the treatment of longitudinal two-stream instabilities and their relation to the longitudinal impedance will be addressed in a future study.

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