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Effect of sheath potential on electromagnetic radiation emitted from the rear surface of a metallic foil target*

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In ultra-intense laser–matter interactions, intense electric fields formed at the rear surface of a foil target may have strong influences on the motion of energetic electrons, and thereby affect the electromagnetic emissions from the rear surface, usually ascribed to transition radiation. Due to the electric fields, transition radiation occurs twice and bremsstrahlung radiation also happens because the electrons will cross the rear surface twice and have large accelerations. In the optic region, transition radiation is dominant. The radiation spectrum depends on the electric field only when the electrons are monochromatic, and becomes independent of the electric field when the electrons have a broadband momentum distribution. Therefore, in an actual experiment, the electric field at the rear surface of a foil could not be studied just with the measurement of optic emissions. In the terahertz region, both bremsstrahlung and transition radiations should be taken into account, and the radiation power could be enhanced in comparison with that without the inclusion of bremsstrahlung radiation. The frequency at which the maximum terahertz radiation appears depends on the electric field.

Keywords: sheath potential, transition radiation, bremsstrahlung radiation, energetic electrons
PACC: 5225P, 5270K

1. Introduction

Optic emissions from rear surface of a metallic foil irradiated with an ultra-intense laser pulse have been extensively investigated in recent years.[1−7] Energetic electron beams generated in ultra-intense laser–matter interactions can be characterized through the measurements of the optic emissions from the rear surface of a foil target. Important aspects of the beams have been inferred, such as beam transport,[1,6] beam bunches,[2] and beam generations,[3,4] etc. Lower frequency radiation, i.e., terahertz emission has also been observed from a laser-driven electron beam crossing a plasma-vacuum boundary.[8] Nowadays, the optic emissions from the rear surface of a metallic foil are mainly ascribed to optic transition radiations that are emitted when energetic electrons cross the interface between two different media,[9] i.e. the vacuum and the foil. Based upon this assumption, optic emissions from rear surface of a metallic foil and terahertz radiation from plasma-vacuum boundary are also theoretically investigated.[10−12] Usually, uniformly moving charges are assumed in calculations.

In reality, however, extreme quasi-static electric fields can appear around the rear surface of a foil when energetic electrons generated at the front surface propagate through the target and form a dense electron plasma sheath at the rear surface. The electric fields in the sheath are so strong that they can ionize the material around the rear surface of the target, and accelerate ions up to multi-MeV energies.[13,14] The observation of the generation of multi-MeV ions is an evidence of the plasma sheath formed by the energetic electrons.[15,16] Very recently, the electric fields in plasma sheath are directly detected with a beam of protons that is generated by an auxiliary ultra-intense laser pulse.[17] From the deflection of the protons, the electric field is estimated as high as 3×10^{10} V/m. Obviously, these electric fields can strongly affect the motions of the energetic electrons: most of the energetic electrons are drawn back to the target by the

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electric fields when they leave the rear surface. In consequence, the intense electric fields at the rear surface of a foil can affect optic emissions at least in two aspects. The first is that transition radiations should occur twice as an electron leaves and goes back to the rear surface. The second is that bremsstrahlung radiations should be emitted due to the accelerations of the electrons in the intense electric field. To the best knowledge of the authors, neither of these two factors is considered in a theory of transition radiation. Therefore, it is necessary to survey the effect of an intense electric field on the electromagnetic emissions from the rear surface of a foil target. In this article, we present a simplified model of the electromagnetic radiations from the rear surface of a metallic foil with the consideration of transition radiation and bremsstrahlung radiation. Our calculations show that in the optic region where bremsstrahlung radiation is negligible the spectra of optic emissions could be influenced by the electric fields. In the lower frequency region, i.e., terahertz region, the radiation intensity could be enhanced due to bremsstrahlung radiation.

2. Model, calculations and discussions

Our calculations are based on a simplified model. As shown in Fig.1, there exists a uniform static electric field $\boldsymbol{E}$ along the normal of a plate target which is made from a perfect conductor. The target normal is parallel to the x-axis. An electron $-$e leaves the target surface with a momentum of p, and backs to the target under the action of the electric field. For simplicity, the electron is assumed to move just along the x-axis.

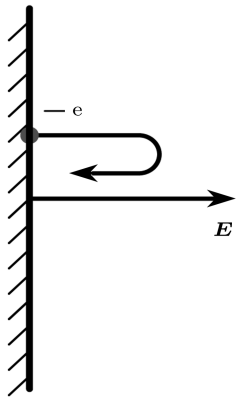


Fig.1. The model of the calculations.

Before leaving or after going back to the surface, the electric fields together with the electrons disappear because of the screening of the perfect conductor. When the electron moves in the vacuum, the effect of the interface between the target and vacuum can be taken into account by including so-called image charge. The image charge has an opposite charge number, and has a velocity anti-parallel to the electron. With this model, the vector potentials $\boldsymbol{A}_{\mp}$ of the electron and its image charge in the wave zone can then be respectively written as

$$\boldsymbol{A}_{-}(\boldsymbol{R},t)=\begin{cases}-\dfrac{e\boldsymbol{v}(t')}{cR(1-\boldsymbol{n}\cdot\boldsymbol{v}(t')/c)}, & 0<t'<2\tau,\\ 0, & t'<0 \quad \text{or} \quad t'>2\tau;\end{cases} \tag{1a}$$

$$\boldsymbol{A}_{+}(\boldsymbol{R},t)=\begin{cases}-\dfrac{e\boldsymbol{v}(t'')}{cR(1+\boldsymbol{n}\cdot\boldsymbol{v}(t'')/c)}, & 0<t''<2\tau,\\ 0, & t''<0 \quad \text{or} \quad t''>2\tau.\end{cases} \tag{1b}$$

Here $\boldsymbol{R}$ is the observation point in the wave zone, $\boldsymbol{n}=\boldsymbol{R}/R$ is the radiation direction, $\boldsymbol{v}(t')$ is the electron velocity, and t' and t'' are the retardation times determined by

$$t'-\frac{1}{c}\boldsymbol{n}\cdot\boldsymbol{r}(t')=t-R/c, \tag{2a}$$

$$t''+\frac{1}{c}\boldsymbol{n}\cdot\boldsymbol{r}(t'')=t-R/c, \tag{2b}$$

and $\boldsymbol{r}(t')$ is the trajectory of the electron. Here we assume that at $t=0$ the electron leaves the target surface, and at $t=2\tau$ the electron returns the target surface. The spectral intensity of the radiation can then be obtained

$$\frac{\mathrm{d}^2\mathcal{E}_{\boldsymbol{n}\omega}}{\mathrm{d}\Omega\mathrm{d}\omega}=\frac{e^2\omega^2}{\pi^2c^3}\left|\int_0^{2\tau}\boldsymbol{v}(t)\times\boldsymbol{n}\mathrm{e}^{\mathrm{i}\omega t}\cos\left[\omega\boldsymbol{n}\cdot\boldsymbol{r}(t)/c\right]\mathrm{d}t\right|^2. \tag{3}$$

In Eq.(3), the lower and upper limits of the integral are finite, and the charge velocity $\boldsymbol{v}(t)$ is also finite at the limits of the integral. It is easy to demonstrate that the integral in Eq.(3) is not zero even though the charge velocity is constant. This is due to transition radiation that occurs when the charge leaves and goes back to the target surface. When the charge is accelerated, its velocity is no longer a constant. Hence bremsstrahlung radiation is naturally included in Eq.(3).

The motion of an electron in a uniform electric field can be easily calculated. The electron velocity and trajectory are given by

$$\boldsymbol{v} = c\frac{(p/m_{\mathrm{e}}c)(1-t/\tau)}{\sqrt{1+(p/m_{\mathrm{e}}c)^2(1-t/\tau)^2}}\boldsymbol{e}_x, \tag{4a}$$

$$\boldsymbol{r} = \frac{m_{\mathrm{e}}c^2}{eE}\left[\sqrt{1+(p/m_{\mathrm{e}}c)^2} - \sqrt{1+(p/m_{\mathrm{e}}c)^2(1-t/\tau)^2}\right]\boldsymbol{e}_x, \tag{4b}$$

where m_{e} is the electron mass, p is the momentum of the electron at $t=0$, $\tau = p/\mathrm{e}E$ is the time when the electron velocity becomes zero, and $\boldsymbol{e}_x$ is the unit vector along the x-axis. Substituting Eq.(4) into Eq.(3) and performing some basic calculations, the spectral intensity of the radiation can be written as,

$$\frac{\mathrm{d}^2\mathcal{E}}{\mathrm{d}\omega\mathrm{d}\Omega} = \frac{e^2\omega^2\tau^2}{\pi^2 c}\sin^2\theta\left|\int_{-1}^{1}\mathrm{d}\xi \mathrm{e}^{\mathrm{i}\omega\tau\xi}\frac{(p/m_{\mathrm{e}}c)\xi}{\sqrt{1+(p/m_{\mathrm{e}}c)^2\xi^2}} \times\cos\left\{\frac{\omega\tau}{p/m_{\mathrm{e}}c}\left[\sqrt{1+\frac{p^2}{m_{\mathrm{e}}^2c^2}}-\sqrt{1+\frac{p^2\xi^2}{m_{\mathrm{e}}^2c^2}}\right]\cos\theta\right\}\right|^2, \tag{5}$$

where θ is the angle between the target normal and the wave vector of the radiation. Introducing a function $G(\omega,p,\theta)$ which is defined as

$$G(\omega,p,\theta) = \omega\int_{-1}^{1}\mathrm{d}\xi\mathrm{e}^{\mathrm{i}\omega\xi}\frac{p\xi}{\sqrt{1+p^2\xi^2}}\cos\left\{\frac{\omega}{p}\left[\sqrt{1+p^2}-\sqrt{1+p^2\xi^2}\right]\cos\theta\right\}^2, \tag{6}$$

the spectral intensity (5) can be simplified as

$$\frac{\mathrm{d}^2\mathcal{E}}{\mathrm{d}\omega\mathrm{d}\Omega} = \frac{e^2}{\pi^2 c}\sin^2\theta\left|G(\omega\tau, p/m_{\mathrm{e}}c,\theta)\right|^2. \tag{7}$$

Hence the spectral intensity depends on the electron momentum and the electric field through function G. We then give a brief discussion on the property of the function G.

The integral on the right hand side of Eq.(6) is rather complicated. We only discuss the function G in several cases of interest. When $\theta=\pi/2$, the function G defined with Eq.(6) can be greatly simplified,

$$G(\omega,p,\theta=\pi/2) = \omega\int_{-1}^{1}\mathrm{d}\xi\mathrm{e}^{\mathrm{i}\omega\xi}\frac{p\xi}{\sqrt{1+p^2\xi^2}}. \tag{8}$$

In the high frequency limit, i.e., $\omega\gg 1$, the asymptotic expansion of $G(\omega,p,\theta=\pi/2)$ is given by,

$$G(\omega,p,\theta=\pi/2) = \mathrm{i}\frac{2p}{\sqrt{1+p^2}}\cos\omega + O(\omega^{-1}). \tag{9}$$

This result shows that $G(\omega,p,\theta=\pi/2)$ is an oscillating function of ω in the high frequency limit. In the low frequency limit, i.e. $\omega\ll 1$, we have

$$G(\omega,p,\theta=\pi/2) = \mathrm{i}2\omega^2\frac{p\sqrt{1+p^2}+\ln(\sqrt{1+p^2}+p)}{p^2} + o(\omega^4). \tag{10}$$

We see that the function $G(\omega,p,\theta=\pi/2)$ decreases to zero very quickly when the normalized frequency ω approaches zero. We numerically integrate Eq.(6), and plot the function $|G(\omega,p,\theta)|^2$ versus ω in Fig.2 at $p=1.7$ and $\theta=\pi/2$. As shown in Fig.2, $|G|^2$ is indeed an oscillating function, and becomes 0 when $\omega=0$. The first maximum is significantly larger than others. When $\omega\gg 1$, the amplitudes of the oscillations approach to the value predicted by Eq.(9).

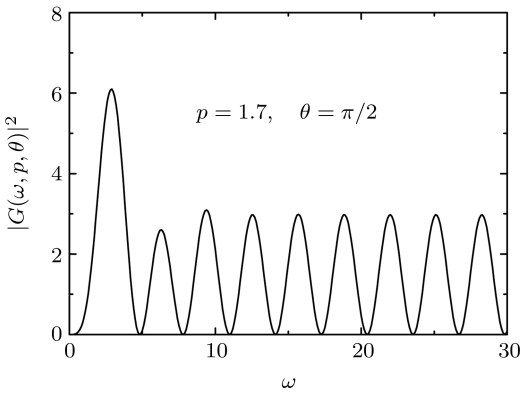


Fig.2. The function $G(\omega, p, \theta)$ versus ω when $p = 1.7$ and $\theta = \pi/2$.

The physics hinted at in Fig.2 is quite understandable. In the high frequency region of $\omega\tau \gg 1$, bremsstrahlung radiation is negligible.[18] Thus the radiation in this region is mainly due to transition radiation. In our model, transition radiation occurs twice when the electron leaves and goes back to the surface. There is a phase difference of $2\omega\tau$ between the radiation fields. An interference term $|1 + \exp(\mathrm{i}2\omega\tau)|^2 \propto \cos^2(\omega\tau)$ then arises in the spectral intensity. This is the physical reason for the asymptotic expansion of Eq.(9). In the low-frequency region, the first maximum of the spectral intensity has the highest amplitude, as shown in Fig.2. It has been known that bremsstrahlung radiation becomes significant in the low-frequency region of $\omega\tau \lesssim 1$ in the case of our interest.[18] Therefore, the radiation becomes more intense in low-frequency region because both transition radiation and bremsstrahlung radiation contribute to radiation energies.

Since the spectral intensity is nearly proportional to the factor $\cos^2 \omega\tau$ in the high-frequency region, which is an oscillating function, the electric field could be inferred if the modulation frequency of the radiation spectrum is measured in case that the initial momentum of the charge is known,

$$\Delta\omega = \frac{\pi eE}{p}. \tag{11}$$

In principle, Eq.(11) provides us a method of the measurement of the electric field.

When a beam of N electrons leaves and goes back to the target, the spectral intensity of the radiation is now given by

$$\frac{\mathrm{d}^2\mathcal{E}}{\mathrm{d}\omega\mathrm{d}\Omega} = \frac{e^2\omega^2}{\pi^2 c^3}\left|\sum_{i=1}^{N}\int_{\tau_i}^{\tau_i+2p_i/eE} \boldsymbol{v}_i(t)\times\boldsymbol{n}\mathrm{e}^{-\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{\rho}_i}\cos[\boldsymbol{k}\cdot\boldsymbol{r}_i(t)]\mathrm{e}^{\mathrm{i}\omega t}\mathrm{d}t\right|^2,$$

where $\boldsymbol{k} = \omega\boldsymbol{n}/c$ is the wave vector of the radiation, τ_i is the time when the i-th electron leaves the surface, $\boldsymbol{r}_i(t) = \boldsymbol{\rho}_i + x_i(t)\boldsymbol{e}_x$ is the trajectory of the i-th electron, and $\boldsymbol{\rho}_i$ is the coordinate where the i-th electron is emitted. The particle velocity and trajectory are given by

$$\boldsymbol{v}_i(t) = c\frac{(p_i/m_\mathrm{e}c)[1-(t-\tau_i)/(p_i/eE)]}{\sqrt{1+(p_i/m_\mathrm{e}c)^2[1-(t-\tau_i)/(p_ieE)]^2}}\boldsymbol{e}_x,$$

$$\boldsymbol{r}_i(t) = \boldsymbol{\rho}_i + c\frac{p_i/eE}{p_i/m_\mathrm{e}c}\left[\sqrt{1+(p_i/m_\mathrm{e}c)^2} - \sqrt{1+(p_i/m_\mathrm{e}c)^2[1-(t-\tau_i)/(p_i/eE)]^2}\right]\boldsymbol{e}_x.$$

Introducing the distribution function

$$F(\tau, \boldsymbol{\rho}, p) = \frac{1}{N}\left\langle\sum_{i=1}^{N}\delta(\tau-\tau_i)\delta(\boldsymbol{\rho}-\boldsymbol{\rho}_i)\delta(p-p_i)\right\rangle, \tag{12}$$

where $\langle\cdots\rangle$ denotes ensemble average, we rewrite the spectral intensity as

$$\begin{aligned}\frac{\mathrm{d}^2\mathcal{E}}{\mathrm{d}\omega\mathrm{d}\Omega} &= N\frac{e^2}{\pi^2 c}\sin^2\theta\int\mathrm{d}\tau\mathrm{d}^2\rho\mathrm{d}pF(\tau,\boldsymbol{\rho},p)\,|G(\omega p/eE, p/m_\mathrm{e}c, \theta)|^2 \\ &\quad + N(N-1)\frac{e^2}{\pi^2 c}\sin^2\theta\left|\int\mathrm{d}\tau\mathrm{d}^2\rho\mathrm{d}pF(\tau,\boldsymbol{\rho},p)G(\omega p/eE, p/m_\mathrm{e}c,\theta)\mathrm{e}^{\mathrm{i}\omega(\tau+p/eE)}\mathrm{e}^{-\mathrm{i}\boldsymbol{k}.\boldsymbol{\rho}}\right|^2. \end{aligned}\tag{13}$$

The first term on the right hand side of Eq.(13) is proportional to the total number of the electrons in the beam. This term, usually called incoherent term also depends on the momentum distribution of the electrons.

The momentum distribution function of the electrons in the beam is given by

$$F_p(p) = \int \mathrm{d}\tau \mathrm{d}^2\rho F(\tau, \boldsymbol{\rho}, p). \tag{14}$$

With the momentum distribution function $F_p(p)$, the incoherent term of the spectral intensity (13) can be written as

$$\frac{\mathrm{d}^2\mathcal{E}_I}{\mathrm{d}\omega \mathrm{d}\Omega} = N\frac{e^2}{\pi^2 c}\sin^2\theta \int \mathrm{d}p F_p(p) \left|G(\omega p/eE, p/m_\mathrm{e}c, \theta)\right|^2. \tag{15}$$

The second term on the right-hand side of Eq.(13) is proportional to the square of the electron number when $N \gg 1$. This term, called coherent term, comes from the interferences of the radiations emitted by different electrons. It is a functional of the bunch form of the beam, and can be written as

$$\frac{\mathrm{d}^2\mathcal{E}_\mathrm{C}}{\mathrm{d}\omega \mathrm{d}\Omega} = N(N-1)\frac{e^2}{\pi^2 c}\sin^2\theta \left|\int \mathrm{d}p \tilde{F}(\omega, \boldsymbol{k}_\perp, p) G(\omega p/eE, p/m_\mathrm{e}c, \theta)\mathrm{e}^{\mathrm{i}\omega p/eE}\right|^2, \tag{16}$$

where $\tilde{F}(\omega, \boldsymbol{k}_\perp, p)$ is the Fourier component of the distribution function,

$$\tilde{F}(\omega, \boldsymbol{k}_\perp, p) = \int \mathrm{d}\tau \mathrm{d}^2\rho F(\tau, \boldsymbol{\rho}, p)\mathrm{e}^{\mathrm{i}\omega\tau - \mathrm{i}\boldsymbol{k}_\perp\cdot\boldsymbol{\rho}}. \tag{17}$$

We first discuss the coherent term in the optic region because coherent radiation is much more intense than incoherent radiation, and is more easily detected in experiments.[2,3] For an experiment of interest, we take $E = 3\times10^{10}\,\mathrm{V/m}$ and $p/m_\mathrm{e}c = 1.7$. Then we have $p/eE = 9.7\times10^{-14}\,\mathrm{s}$. In the optic region, say around the wavelength of $0.532\,\mu\mathrm{m}$, the frequency ω is about $3.5\times10^{15}\,\mathrm{Hz}$. Then we have $\omega p/eE \gg 1$. In this case, the coherent term (16) can be approximately written as

$$\frac{\mathrm{d}^2\mathcal{E}_\mathrm{C}}{\mathrm{d}\omega \mathrm{d}\Omega} \approx 4N(N-1)\frac{e^2}{\pi^2 c}\sin^2\theta \left|\int \mathrm{d}p \tilde{F}(\omega, \boldsymbol{k}_\perp, p)\frac{p/m_\mathrm{e}c}{\sqrt{1+(p/m_\mathrm{e}c)^2}}\cos(\omega p/eE)\,\mathrm{e}^{\mathrm{i}\omega p/eE}\right|^2. \tag{18}$$

If the electron beam is monochromatic, i.e. $\tilde{F}(\omega, \boldsymbol{k}_\perp, p) = n(\omega, \boldsymbol{k}_\perp)\delta(p - p_0)$, where $n(\omega, \boldsymbol{k}_\perp)$ is the Fourier component of the beam bunch form, then we have

$$\frac{\mathrm{d}^2\mathcal{E}_\mathrm{C}}{\mathrm{d}\omega \mathrm{d}\Omega} \approx 4N(N-1)\frac{e^2}{\pi^2 c}\frac{p_0^2}{m_\mathrm{e}^2c^2 + p_0^2}\sin^2\theta\cos^2(\omega p_0/eE)\left|n(\omega, \boldsymbol{k}_\perp)\right|^2. \tag{19}$$

In this case, the spectral intensity not only depends on the bunch form of the beam through the factor $|n(\omega, \boldsymbol{k}_\perp)|^2$ but also depends on the electric field through the factor $\cos^2(\omega p_0/eE)$. If the beam has a Gaussian bunch form, the spectral intensity of the radiation then has a Gaussian envelope modulated with the factor $\cos^2(\omega p_0/eE)$. We can then infer the electric field from the radiation spectrum using Eq.(11). Monochromatic electrons are usually generated with gas jet target.[19−21] In laser-solid interaction, energetic electrons usually have a Boltzmann energy distribution.[3,22] The momentum distribution function of the electrons is broadband. In this case, the integration over the fast oscillating factor becomes negligible. Then we have approximately

$$\frac{\mathrm{d}^2\mathcal{E}_\mathrm{C}}{\mathrm{d}\omega \mathrm{d}\Omega} \approx N(N-1)\frac{e^2}{\pi^2 c}\sin^2\theta \left|\int \mathrm{d}p \tilde{F}(\omega, \boldsymbol{k}_\perp, p)\frac{p/m_\mathrm{e}c}{\sqrt{1+(p/m_\mathrm{e}c)^2}}\right|^2. \tag{20}$$

The coherent term is now independent of the electric field. In an actual experiment, the spectral intensity of the radiation in the optic region should lie between the two limiting cases described respectively by Eqs.(19) and (20). However, the modulation of the radiation spectrum should be very small so as not to be visible. Hence it seems impossible that the electric field could be derived from the spectra intensity of the optic

emissions.

The temperature of energetic electrons can also affect the radiation intensity. In the case that $\tilde{F}(\omega, \boldsymbol{k}_\perp, p)$ is separable, $\tilde{F}(\omega, \boldsymbol{k}_\perp, p) = n(\omega, \boldsymbol{k}_\perp) \exp[-(m_e^2c^4 + p^2c^2)^{1/2}/T]$, where T is the temperature of the hot electrons, we have

$$\frac{\mathrm{d}^2\mathcal{E}_\mathrm{C}}{\mathrm{d}\omega \mathrm{d}\Omega} \approx N(N-1)\frac{e^2}{\pi^2 c}\frac{T^2}{m_e^2c^4}\sin^2\theta |n(\omega, \boldsymbol{k}_\perp)|^2.$$

Here we can see that the radiation energy is proportional to the square of the hot electron temperature.

Radiation in terahertz region can also be emitted as the electron beam cross the rear surface of the target. Powerful terahertz radiation has been observed from a laser-driven electron bunch crossing a plasma–vacuum boundary by using a gas target.[8] However, the effect of electric fields is not included in the calculations.[12] As clearly shown in Fig.2, the radiation could be strongly affected by the electric field. Especially, the radiation intensity could be enhanced due to the contribution of bremsstrahlung. For an experiment of interest, we take $E = 3\times10^{10}\,\mathrm{V/m}$[17] and $p/m_ec = 1.7$.[3] As shown in Fig.2, the first maximum of the radiation appears at the normalized frequency of $\omega p/eE \sim 3$, i.e., $\omega \sim 3\times10^{13}\,\mathrm{Hz}$, which is just in terahertz region. This result shows that electric fields may play a significant role in terahertz emission from laser-driven electron beam crossing the rear surface of a foil. On the other hand, the electric field may be inferred from terahertz radiation. Since the maximum radiation appears around $\omega p/eE \sim 3$, the electric field may be obtained from the frequency corresponding to the maximum terahertz intensity.

3. Summary

In summary, based upon a simplified model, we have investigated the effect of the electric field in a potential sheath on the electromagnetic radiation emitted from the rear surface of a foil. Due to the intense electric field, most energetic electrons generated at the front surface cannot leave the target. They are attracted by the electric field, and back to the target. As the electrons leave and go back to the rear surface of the target, both bremsstrahlung and transition radiations are emitted. In the optic region, transition radiation is dominant. The spectral intensity of coherent transition radiation is hardly affected by the presence of electric field except in the case that the energetic electrons are monochromatic. Hence the strength of electric field cannot be inferred from the optic spectrum emitted from the rear surface of the target. The radiation in the terahertz region is strongly affected by the presence of the electric field. Terahertz radiation may be enhanced due to the electric field. And the electric field may be inferred from the frequency at which the maximum terahertz radiation appears. It should be pointed out that our model is only suitable for an experiment with metallic foil. For a transparent target experiment, Cherenkov radiation from the energetic electrons becomes dominant, which has been investigated both in experiment[6] and theory.[23]

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