Radiative transfer dynamo effect

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Magnetic fields in rotating and radiating astrophysical plasma can be produced due to a radiative interaction between plasma layers moving relative to each other. The efficiency of current drive, and with it the associated dynamo effect, is considered in a number of limits. It is shown here, however, that predictions for these generated magnetic fields can be significantly higher when kinetic effects, previously neglected, are taken into account.

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I. INTRODUCTION

Cosmic magnetism is usually explained by magnetohydrodynamic dynamo theory [1], which is, however, only an amplification mechanism that still requires some initial seed field. There have been many speculations about the origin of the seed field, but consensus is still lacking [2,3]. One possible mechanism is a radiation induced drag force on electrons in rotating astrophysical objects. This idea was evidently first proposed by Cattani and Sacchi [4] and later has been applied to different astrophysical conditions and objects [5–17]. However, none of these studies took into account kinetic effects. It is shown here that predictions for these generated magnetic fields can be significantly higher when kinetic effects are taken into account. In the presence of existing magnetic fields, these kinetic effects can enhance the generated magnetic fields by orders of magnitude.

A rotating astrophysical object is subject to asymmetric incoming radiation, which exerts the Poynting-Robertson drag force on electrons in the (toroidal) rotation direction that leads to the poloidal magnetic field. Within a fluid framework, this can be modeled by including an additional term into the equation for the magnetic field dynamics:

$$\frac{\partial \mathbf{B}}{\partial t} = -\frac{c}{e}\nabla \times \mathbf{f}_{\mathrm{rad}}. \tag{1}$$

There are two ways in which kinetic effects modify the effective radiation force. First, the Poynting-Robertson force on an individual electron depends on the absorbed power, which is, generally speaking, different for the electrons of different energies; usually the more energetic electrons absorb more power. Thus, to get the effective radiation force on the electron fluid, one needs to average the force for each electron over the absorbed power. Second, toroidal current can be driven even without toroidal momentum injection just by asymmetrically heating electrons. Indeed, by heating electrons we increase their energy and since the collision frequency in plasma is energy dependent ($\propto v^{-3}$) the toroidal drag force due to Coulomb collisions is going to be asymmetric resulting in the total toroidal current [18].

To simplify the problem and underscore the influence of the kinetic effects, we consider a slab geometry, where the parallel direction corresponds to the toroidal direction of the original rotating object (see Fig. 1). Namely, we consider two parallel and possibly magnetized (in the parallel direction) plasma slabs that move relative to each other with velocity $\bar{\beta}$ (velocities are measured in the units of c). We label the upper slab slab 1 and the lower slab slab 2.

The paper is organized as follows: In Sec. II we derive the efficiency of current generation through the Poynting-Robertson effect. In Sec. III, using kinetic rather than fluid theory, we show how the efficiency of the current generation through radiation effects can be much enhanced when there is a seed magnetic field already present and when kinetic effects are considered. We consider, in Sec. III A, the case of blackbody emission and cyclotron absorption. In Sec. III B, we consider the case of cyclotron emission and cyclotron absorption, where not only can the currents driven be driven much more effectively, but there is even the curious effect that the current in adjacent differentially moving plasma can be either in the same direction or in opposite directions. In Sec. IV we summarize and discuss our findings.

II. THE POYNTING-ROBERTSON EFFECT

Consider an electron that moves with velocity $\beta_\parallel$ and emits isotropic radiation in its own reference frame. Imagine that this electron also absorbs external radiation, which is isotropic in its own reference frame moving with parallel velocity $\beta_s = -\bar{\beta}$. Conservation of energy and momentum then gives

$$mc(\gamma\dot{\beta}_\parallel + \dot{\gamma}\beta_\parallel) + \dot{p}_\parallel^{\mathrm{ems}} = \dot{p}_\parallel^{\mathrm{abs}}, \tag{2}$$

$$mc^2\dot{\gamma} = P^{\mathrm{abs}} - P^{\mathrm{ems}}, \tag{3}$$

where P^{abs} is the absorbed power, $p_\parallel^{\mathrm{abs}}$ is the absorbed parallel momentum, P^{ems} is the emitted power, and $p_\parallel^{\mathrm{ems}}$ is the emitted parallel momentum.

The time derivative of the wave momentum is determined by the power delivered by the wave $\dot{\mathbf{p}}^{\mathrm{wave}} = (\mathbf{k}/\omega)P^{\mathrm{wave}}$. Using the Lorentz transformation we can express it as $\dot{p}^{\mathrm{ems}} = (\beta_\parallel/c)P^{\mathrm{ems}}, \dot{p}^{\mathrm{abs}} = (\beta_s/c)P^{\mathrm{abs}}$. Inserting these expressions into the energy-momentum equations we find that electron parallel velocity satisfies

$$\dot{\beta}_\parallel = -\frac{P^{\mathrm{abs}}}{\gamma mc^2}(\beta_\parallel - \beta_s). \tag{4}$$

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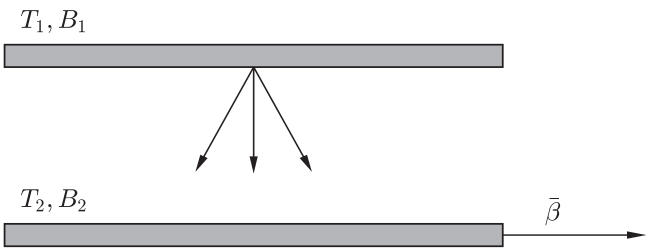

FIG. 1. Parallel radiating and absorbing slabs of plasma, immersed in different magnetic fields at different temperatures, in relative parallel motion.

We see that the electron experience drag by absorbing the external radiation. This effect is called the Poynting Robertson effect. It is a relativistic effect by its very nature, although it does not require that the relative velocity between absorber and emitter be relativistic. In the reference frame of an electron, this drag force can be simply interpreted as a momentum transfer from asymmetric external radiation to an electron. However, in the reference frame of the external radiation source, there is no parallel momentum injection; there is only an energy increase, which makes the electron heavier relativistically. Since the total parallel momentum must be conserved, the electron must slow down. Notice that, without external radiation, there would be no radiation drag force (if we define force as the cause of velocity change rather than momentum), which is consistent with the well-known fact that isotropically radiating charge conserves its parallel velocity [19]. It should be emphasized that here, by absorption, we mean a generalized process of wave-particle interaction; for example, it can denote Thomson scattering. While electrons experience radiation drag, ions are almost unaffected by radiation and hence current is generated.

Let us estimate crudely the current drive efficiency in this case. Assume that the parallel velocity is randomized during the inverse collision time ν^{-1} and use the effective electron-electron and electron-proton Coulomb collision frequency $\nu = \Gamma/6\beta^3$ (see Refs. [20–22]), where $\Gamma = \omega_p^4 \ln \Lambda / 4\pi n c^3$. Then, after averaging over the Maxwell distribution, we find

$$\frac{j_\parallel}{P_V^{\text{abs}}} = \frac{\langle \beta^3 \rangle}{15} \bar{\beta} \approx 0.43 \beta_{\text{th}}^3 \bar{\beta}. \tag{5}$$

Here and later the current drive efficiency is expressed in the units of $e/\Gamma mc$ except for Eq. (6).

III. KINETIC FORMULATION

The kinetic theory of current drive by external radiation is well developed and experimentally demonstrated [23]. This theory has been advanced to accommodate the need for efficient noninductive toroidal current generation required for the successful operation of commercial tokamaks. The theory formulates the efficiency of current generation as the ratio of the driven current density to the absorbed power density [22]:

$$\frac{j_\parallel}{P_V^{\text{abs}}} = -\frac{e}{mc}\left[\frac{n_\parallel}{\nu} + \frac{\beta_\parallel}{\beta}\frac{\partial}{\partial \beta}\left(\frac{1}{\nu}\right)\right]. \tag{6}$$

The first term in Eq. (6) can be associated with the Poynting-Robertson drag, while the second is due to asymmetric heating. The second term arises because the collision frequency ν depends sensitively on the electron energy. It leads to the electron cyclotron current drive effect in tokamaks [18]. The radiative transfer dynamo effect in astrophysical contexts is not much different from the situation described above. The major difference is that the radiation driving current is set up naturally rather than controlled.

Equation (6) gives the nonrelativistic efficiency of the current driven by a very narrow radiation band that affects only a small region in velocity space. To calculate the efficiency for arbitrary incoming radiation average Eq. (6) over the power density absorbed per frequency per solid angle per $d^3\boldsymbol{\beta}$. The absorbed power density is given by

$$P_V^{\text{abs}} = \iint d\omega\, d\Omega \alpha_{\omega\Omega} I_{\omega\Omega}, \tag{7}$$

where $I_{\omega\Omega}$ is the incoming electromagnetic energy flux density per unit frequency per solid angle and $\alpha_{\omega\Omega}$ is the absorption coefficient (true absorption plus stimulated emission). Due to the principle of detailed balance the absorption coefficient can be expressed through the emissivity of an individual electron $\eta_{\omega\Omega}(\mathbf{p})$ [24,25]:

$$\alpha_{\omega\Omega} = -\frac{8\pi^3 c^2}{n_r^2 \omega^2} \int d^3\mathbf{p}\, \eta_{\omega\Omega}(\mathbf{p}) \left(\frac{\partial f}{\partial \varepsilon} + \frac{n_\parallel}{c}\frac{\partial f}{\partial p_\parallel}\right), \tag{8}$$

where n_r is the ray refractive index, we will use approximation of tenuous plasma and assume $n_r \approx 1$; $n_\parallel$ is the wave parallel refractive index, which we take to be just $n_\parallel = \cos\theta$.

Thus, we can calculate the current drive efficiency for a specific type of the absorption mechanism determined by $\eta_{\omega\Omega}(\mathbf{p})$ and for a given external radiation spectrum $I_{\omega\Omega}$. Although, in both emitting and absorbing radiation, the two slabs form a coupled system, to get the efficiency linear in power transferred, note that each slab may be considered to see fixed radiation from the other slab.

We first argue that it is the current drive efficiency that determines large-scale magnetic field generation in optically thick plasma, for which the effect is maximized. For optically thick plasma, the incoming radiation flux $I = \iint d\omega\, d\Omega I_{\omega\Omega}$ is absorbed over the characteristic distance $R = \alpha^{-1}$, where $\alpha = P_V^{\text{abs}}/I$ is the characteristic absorption coefficient. Ampere's law gives $B \cdot 2\pi r \approx (4\pi/c) j_\parallel R h$, where h is the height of the plasma disk, so the large-scale equilibrium magnetic field at distance r outside the plasma is proportional to the current drive efficiency:

$$B \approx \frac{2}{c}\frac{h}{r}\left(\frac{j_\parallel}{P_V^{\text{abs}}}\right) I. \tag{9}$$

Kinetic effects change the current drive efficiency and thus the equilibrium field, but they do not affect much the time to reach equilibrium t_{eq}, which is the so-called "L/R time." That time is still determined by the Spitzer conductivity, since the full distribution function is equally pushed by an electric field [26]. The time to reach equilibrium t_{eq} greatly exceeds the age of the universe, and so the actual value of the magnetic seed is determined at some characteristic time

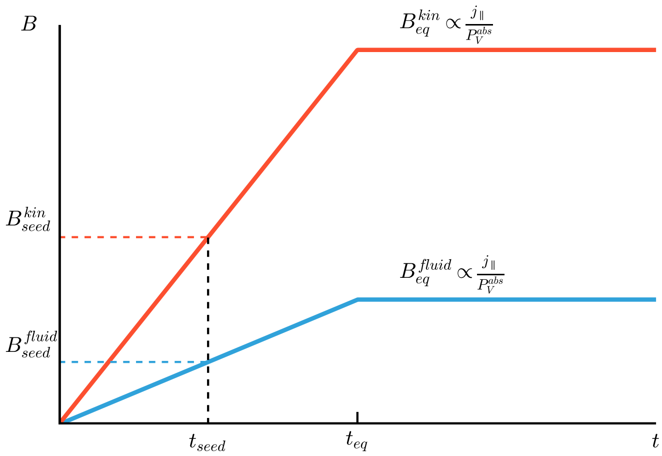

FIG. 2. Schematic picture of the magnetic field growth. Magnetic field grows approximately linearly with time until it saturates at equilibrium value $B_{\rm eq}$. Kinetic effects increase $B_{\rm eq}$ to the same extent as they increase the current drive efficiency, but they hardly change the time to reach equilibrium $t_{\rm eq}$. The actual value of the magnetic seed is determined at some characteristic time $t_{\rm seed} \ll t_{\rm eq}$, and it increases to the same extent as the efficiency increases.

$t_{\rm seed} \ll t_{\rm eq}$, and it increases to the same extent as the efficiency increases (see Fig. 2).

A. Blackbody incoming radiation and cyclotron absorption

To take one example, let us assume that the incoming radiation from the first slab is blackbody:

$$I_{\omega\Omega} = \frac{\omega^2}{8\pi^3 c^2} \frac{T_1}{\bar{\gamma}(1+\bar{\beta}\cos\theta)}. \tag{10}$$

If the plasma were already immersed in a parallel magnetic field, then one of the absorption mechanisms would be synchrotron absorption. In the nonrelativistic case, it is determined by the emissivity [27]:

$$\eta_{\omega\Omega}(\boldsymbol{\beta}) = \frac{e^2\beta_\perp^2\omega^2}{4\pi c}(1+\cos^2\theta)\delta[\omega_{c2} - \omega(1-\beta_\parallel\cos\theta)]. \tag{11}$$

After some algebra it is easy to show that the current drive efficiency in this case is

$$\frac{j_\parallel}{P_V^{\rm abs}} = -\frac{\langle \beta_\perp^2\beta^3 I_2(\beta_\parallel)\rangle + 3\langle \beta_\perp^2 \beta\beta_\parallel I_1(\beta_\parallel)\rangle}{6\langle \beta_\perp^2 I_1(\beta_\parallel)\rangle}, \tag{12}$$

where the averaging is over the initial distribution that is taken to be a Maxwellian, and the following integrals are introduced:

$$I_1(\beta_\parallel) = \int_{-1}^{1} dx \frac{1+x^2}{(1-\beta_\parallel x)^3(1+\bar{\beta}x)}, \tag{13}$$

$$I_2(\beta_\parallel) = \int_{-1}^{1} dx \frac{x(1+x^2)}{(1-\beta_\parallel x)^3(1+\bar{\beta}x)}. \tag{14}$$

Keeping only the terms of the order $O(\beta_\parallel^2), O(\bar{\beta}^2), O(\beta_\parallel\bar{\beta})$ we find

$$\frac{j_\parallel}{P_V^{\rm abs}} = \frac{\langle\beta_\perp^2\beta^3\rangle + 9\langle\beta_\perp^2\beta\beta_\parallel^2\rangle}{15\langle\beta_\perp^2\rangle}\bar{\beta} \approx 2.4\beta_{\rm th}^3\bar{\beta}. \tag{15}$$

If we ignored the second term in the numerator of Eq. (15) and also did not account for $\beta_\perp^2$ in the absorption, then the efficiency would be given by Eq. (5), i.e., correspond to the case of Thomson scattering analyzed through fluid theory.

From comparison of Eq. (15) and Eq. (5), we see that for cyclotron absorption the inclusion of the kinetic effects boosts the generated current by a factor of 6. This is not a huge change, though it is not insignificant either. For reference, the fluid estimates for the galactic magnetic field are about 10^{-19} G [2], while the estimates for the required lower bound on the seed galactic field is about 10^{-14} G [28].

Cyclotron absorption mechanism needs some parallel (toroidal) magnetic field to be already present to work. However, we see that the efficiency (15) is independent of the magnetic field, so it seems that we can get poloidal magnetic field from a very small toroidal field, generated, for example, by the Biermann battery effect [29]. This works only when all the incoming radiation is absorbed within the plasma, so that the effective absorption length is less than the characteristic size of the system. For blackbody incoming radiation flux and cyclotron absorption the effective absorption coefficient is

$$\alpha = \frac{4}{3\pi}\frac{k_B^4}{c^3\sigma_{SB}T_1^3}\omega_{p2}^2\omega_{c2}^2, \tag{16}$$

or $\alpha \approx 10^{-20+n+2b-3k}$ cm^{-1} for $T_1/k_B = 10^k$ K, $n_2 = 10^n$ cm^{-3}, and $B_2 = 10^b$ G. If we take typical protogalactic values $T_1/k_B = 10^4$ K and $n = 1$ cm^{-3}, then for $B_2 = 10^{-20}$ G that realistically can be produced by the Biermann battery the effective absorption length becomes $R \sim 10^{71}$ cm, which is much larger than characteristic size of the system. Thus, the cyclotron absorption mechanism cannot be responsible for the generation of the galactic seed field. However, it might be a very effective mechanism of poloidal magnetic field generation in already highly magnetized objects such as neutron stars.

B. Cyclotron incoming radiation and absorption

So far we considered that the incoming radiation is blackbody. We can expect that if the incoming radiation were narrower in $k_\parallel$, then its absorption would be more asymmetric in parallel velocity of the second slab, which would result in enhanced efficiency.

To investigate this, consider the case where each of the slabs is immersed in an axial magnetic field, though the respective magnetic fields are not not necessarily of the same strength. Suppose that cyclotron radiation is emitted by an optically thin surface layer of depth L. The incoming flux is then given by [27]

$$I_{\omega\Omega} = \frac{n_1 L e^2 \beta_{\rm th1}}{4\pi\sqrt{2\pi}c}\frac{\omega}{|\cos\theta|}(1+\cos^2\theta)e^{-\frac{\left(\frac{\omega-\omega_{c1}}{\omega\cos\theta}+\bar{\beta}\right)^2}{2\beta_{\rm th1}^2}}. \tag{17}$$

The current drive efficiency for cyclotron absorption has the same form as Eq. (12), but with the following definition of I_1, I_2:

$$I_1(\beta_\parallel) = \left(\int_{-\infty}^{-|a|} dx + \int_{|a|}^{\infty} dx\right) \frac{1}{|x|^3} \frac{(x^2+a^2)^2}{(x-a\beta_\parallel)^2} e^{-\frac{(x+b)^2}{2}}, \tag{18}$$

$$I_2(\beta_\parallel) = \left(\int_{-\infty}^{-|a|} dx + \int_{|a|}^{\infty} dx\right) \frac{a}{x^4} \frac{(x^2+a^2)^2}{(x-a\beta_\parallel)^2} e^{-\frac{(x+b)^2}{2}}, \tag{19}$$

where $a = (\omega_{c2} - \omega_{c1})/\omega_{c2}\beta_{\text{th}1} \equiv \triangle\omega_c/\omega_{c2}\beta_{\text{th}1}$, and $b = (\omega_{c1}/\omega_{c2})(\beta_\parallel/\beta_{\text{th}1}) + \bar{\beta}/\beta_{\text{th}1}$.

The first term in the denominator of Eq. (12) with I_1 and I_2 defined above is due to direct parallel momentum injection and so depends on the sign of $\triangle\omega_c$, while the second is due to asymmetric heating. Since now absorption is localized in velocity space and most of the absorbed power goes into heating rather than giving a parallel push, the second term completely dominates, and the efficiency becomes essentially independent of the sign of $\triangle\omega_c$. There are two qualitatively different cases that produce current of different sign: $|a| \ll 1$ (positive current) and $|a| \gtrsim 1$ (negative current). From numerical treatment it appears that for wide range of parameters the efficiency is approximately given by

$$\frac{|j_\parallel|}{P_V^{\text{abs}}} \sim 10^2 \beta_{\text{th}}^2 \bar{\beta}. \tag{20}$$

This is $10^2/\beta_{\text{th}}$ larger efficiency than that for the blackbody radiation, for $T/k_B \approx 10^4$ K is about $\sim 10^5$. Therefore, at least in the case when the plasma already possesses some toroidal magnetic field, one can expect the generated poloidal magnetic field to be orders of magnitude larger than the previous estimates based on the fluid theory.

These results can be understood from the following qualitative picture. The nonrelativistic cyclotron resonance condition for an electron of slab 2 moving with velocity $\beta_{\parallel 2}$ to absorb the radiation emitted by an electron of slab 1 moving with velocity $\beta_{\parallel 1}$ is

$$\omega_{c1} - k_\parallel c(\beta_{\parallel 2} - \beta_{\parallel 1}) = \omega_{c2}. \tag{21}$$

Here we use $k_\parallel$ corresponding to the reference frame where the electron of slab 1 is stationary, and velocities $\beta_{\parallel 1}$, $\beta_{\parallel 2}$ are measured in the frame where slab 2 is stationary.

If $|a| \ll 1$, then essentially $\omega_{c1} \simeq \omega_{c2}$ and the resonance condition is $k_\parallel = 0$ or $\beta_{\parallel 2} = \beta_{\parallel 1}$. The former condition does not depend on the velocities and cannot lead to the asymmetry, the latter condition leads to an asymmetric absorption. Indeed, the electrons with positive parallel velocity $\beta_{\parallel 2} \approx \bar{\beta}$ absorb less power than the electrons with negative parallel velocity $\beta_{\parallel 2} \approx -\bar{\beta}$, because the latter are in resonance with a much larger number of electrons in slab 1. Thus the electrons with negative parallel velocities will experience less Coulomb drag force than the electrons with positive velocities resulting in positive current. This situation is shown in Fig. 3(a).

If $|a| \gtrsim 1$, then the magnetic fields are different $\omega_{c1} \neq \omega_{c2}$ and the resonance condition (21) implies that the electrons of slab 2 with velocity $\beta_{\parallel 2}$ will resonantly interact with the electrons of slab 1 with parallel velocities satisfying

$$\begin{cases} \beta_{\parallel 1} < \beta_{\parallel 2} - |\triangle\omega_c|/\omega_{c1}, \\ \beta_{\parallel 1} > \beta_{\parallel 2} + |\triangle\omega_c|/\omega_{c1}. \end{cases} \tag{22}$$

Thus there is a window in the absorption for each electron. The electrons of slab 2 with negative parallel velocity around $-\bar{\beta}$ will have larger number of electrons of slab 1 in this window and thus will receive less power than the electrons of slab 2 around $\bar{\beta}$. The result is negative current. This situation is shown in Fig. 3(b).

Notice that one gets the same efficiency but with different sign for the current driven in the first slab. For blackbody incoming radiation, this current would be always in the opposite direction, but, interestingly, for cyclotron radiation, it is possible to have the situation when currents in both slabs flow in the same direction. Indeed, since the parameter a depends on temperature it can have different values

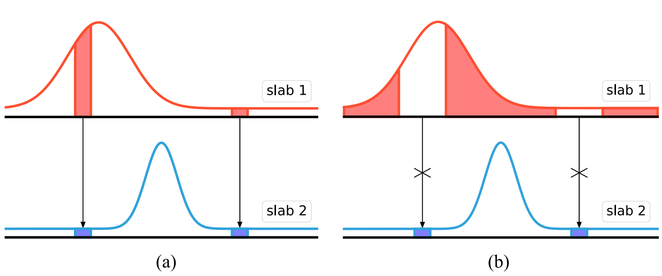

FIG. 3. (a) $\omega_{c1} \simeq \omega_{c2}$: electrons of slab 2 with negative velocities around $-\bar{\beta}$ (left blue region) interact with the large number of electrons of slab 1 (left red region), while symmetric electrons of slab 2 with positive velocities around $\bar{\beta}$ (right blue region) interact with small number of electrons of slab 1 (right red region). (b) $\omega_{c1} \neq \omega_{c2}$: electrons of slab 2 with negative velocities around $-\bar{\beta}$ (left blue region) has large number of electrons of slab 1 in the nonabsorption window (left white region) and thus absorb less energy than symmetric electrons of slab 2 with positive velocities around $\bar{\beta}$ (right blue region), which have small number of electrons of slab 1 in the nonabsorption window (right white region).

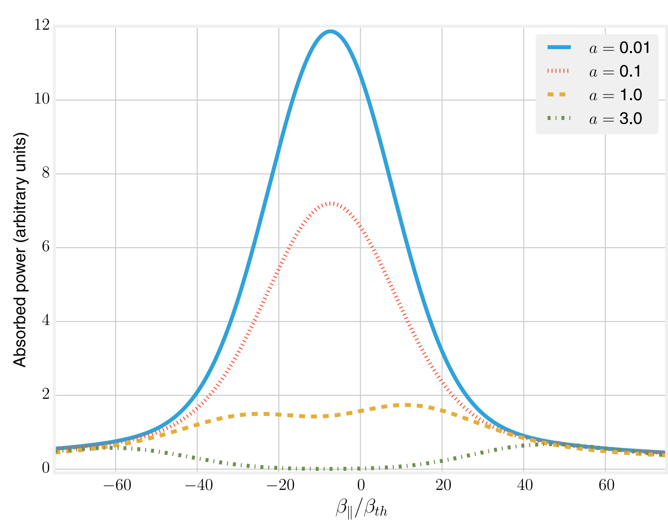

FIG. 4. Absorbed power density per electron as a function of parallel velocity for $\beta_{\text{th}} = 0.01$, $\bar{\beta} = 0.05$ and different values of a: $a = 0.01$ (solid blue), $a = 0.1$ (dotted red), $a = 1.0$ (dashed orange), $a = 3.0$ (dash-dotted green).

corresponding to two different regimes ($|a| \ll 1$ and $|a| \gtrsim 1$) in each slab. Thus, we reach the surprising result that, for a differentially rotating plasma disk immersed in a toroidal magnetic field and with a temperature gradient, it is possible that a toroidal current will be self-consistently generated in the same direction throughout the disk.

Figure 4 shows the absorbed power per electron as a function of the parallel velocity for $\beta_{\rm th} = 0.01$, $\bar{\beta} = 0.05$ and four different values of parameter a. We can clearly see that the absorption is asymmetric. For $|a| \ll 1$ the situation is basically analogous to the case of equal magnetic fields shown in Fig. 3(a) when the electrons with negative parallel velocities (around $-\bar{\beta}$) absorb more power resulting in positive current. In contrast, for $|a| \gtrsim 1$ there is a dip in the absorption for the electrons moving with negative velocities resulting in negative current [see Fig. 3(b)]. We can also see that, as the difference between magnetic fields of the slabs increases, i.e., as the parameter $|a|$ increases, the total absorbed power decreases rapidly. Thus, in the limit $|a| \gg 1$, the efficiency should be viewed as questionable, because the total absorbed power density is negligible and radiation has to pass through a very large volume of plasma to be fully absorbed.

IV. CONCLUSION

The generation of cosmic magnetic fields due to radiation transfer can be significantly larger when one takes into account kinetic effects rather than simply relying on fluid theory. In the case where the radiation is from cyclotron radiation, namely, when there already exists an ambient magnetic field, an increase in fields perpendicular to the ambient field can be orders of magnitude larger when kinetic effects are considered. Curiously, in the case of inhomogeneous field, it is possible to generate these perpendicular fields such that the current produced within two differentially traveling, radiating, and absorbing slabs is in the same direction, an effect that would not be captured in the fluid theory.

The formalism advanced here shows how to deal with a radiative process, which is kinetic by its very nature. It is expected that the formalism advanced here can be applied to various areas of astrophysics where radiative kinetic effects might be important, for example, to radiative magnetic reconnection [30]. It is also hoped that the approach taken here might help to make more accurate the currently widely used astrophysical radiative transfer codes, which, to the best of our knowledge, exist only in the hydrodynamic version (see, for example, Refs. [31,32]).

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