Penetration-depth anisotropy in two-band superconductors

V. G. Kogan and N. V. Zhelezina
Ames Laboratory—DOE and Department of Physics, Iowa State University, Ames, Iowa 50011, USA
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The anisotropy of the London penetration depth is evaluated for two-band superconductors with arbitrary interband and intraband scattering times. If one of the bands is clean and the other is dirty in the absence of interband scattering, the anisotropy is dominated by the Fermi surface of the clean band and is weakly temperature dependent. The interband scattering also suppresses the temperature dependence of the anisotropy.

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The two-gap superconductivity of MgB_2 is established experimentally[1–5] and by solving the Eliashberg equations for the gap distribution on the Fermi surface.[6,7] According to the latter, the gap on the four Fermi surface sheets of this material has two sharp maxima: $\Delta_1 \approx 1.7$ meV at the two π bands and $\Delta_2 \approx 7$ meV at the two σ bands. Within each of these groups, the spread of the gap values is small.

A number of physical properties of MgB_2 were reasonably well described within a model with two constant gaps on two separate Fermi sheets. Still, the data on anisotropy of the magnetic-field penetration depth λ are controversial. The anisotropy parameter $\gamma_\lambda = \lambda_c/\lambda_a$ has been calculated within the *weak-coupling clean-limit* model and shown to increase from about 1.1 at $T=0$ to ≈ 2.6 at T_c.[8] A similar prediction has been made within Eliashberg formalism.[9] Qualitatively, the predictions were confirmed in scanning tunnel microscopy,[10,11] small angle neutron scattering,[12] and magnetization experiments.[13] However, other groups recorded different behavior.[14–17] Given variety of samples used, it is imperative to consider effects of scattering upon γ_λ, a nontrivial problem given different roles of the intraband and interband scattering in two-band materials. The problem has been discussed by Golubov *et al.*,[9] who focused on $\lambda(T)$ in the absence of the interband scattering. The dirty limit λ has been considered by Moca[18] within Eliashberg formalism and by Gurevich,[19] who used the weak-coupling Usadel approach. It is shown below how arbitrary interband and intraband scatterings can be treated within the weak-coupling model.

Our approach is based on the quasiclassical version of the BCS theory for a general anisotropic Fermi surface and for an arbitrary anisotropic order parameter $\Delta(\boldsymbol{k})$.[20] In the absence of currents and fields we have for the Eilenberger Green's functions $f(\boldsymbol{k},\omega)$ and $g(\boldsymbol{k},\omega)$,

$$0 = 2\Delta g - 2\omega f + I, \quad 1 = g^2 + f^2. \tag{1}$$

Here the scattering term I is given by the integral over the full Fermi surface,

$$I(\boldsymbol{k}) = \int d^2\boldsymbol{q}\, \rho(\boldsymbol{q})\, W(\boldsymbol{k},\boldsymbol{q})[g(\boldsymbol{k})f(\boldsymbol{q}) - f(\boldsymbol{k})g(\boldsymbol{q})] \tag{2}$$

with $W(\boldsymbol{k},\boldsymbol{q})$ being the scattering probability from $\boldsymbol{q}$ to $\boldsymbol{k}$ at the Fermi surface. The Matsubara frequencies are $\omega = \pi T(2n+1)$ with an integer n ($\hbar = 1$). The *local* density of state $\rho(\boldsymbol{q})$ is normalized: $\int d^2\boldsymbol{q}\, \rho(\boldsymbol{q}) = 1$.

The system (1),(2) should be complemented with an equation for $\Delta(\boldsymbol{k})$. We will not use it here, rather taking $\Delta(\boldsymbol{k})$ as a given. This simplifies the problem greatly because solving for $\Delta(\boldsymbol{k})$ usually involves a number of assumptions which are difficult to control.

We use the approximation of the scattering time τ,

$$\int d^2\boldsymbol{q}\, \rho(\boldsymbol{q})\, W(\boldsymbol{k},\boldsymbol{q})\, f(\boldsymbol{q}) = \langle f \rangle/\tau; \tag{3}$$

$\langle \cdots \rangle$ stands for the average over the Fermi surface. Clearly, the approximation amounts to the scattering probability $W = 1/\tau$ being constant for any $\boldsymbol{k}$ and $\boldsymbol{q}$.

For two well-separated Fermi surface sheets, the probabilities of intraband scatterings may differ from each other and from processes involving $\boldsymbol{k}$ and $\boldsymbol{q}$ from different bands. The effects of the interband and intraband scattering upon various properties of the system are different, e.g., the intraband scattering does not affect T_c, whereas the interband does. Therefore, Eq. (3) is replaced by

$$\int d^2\boldsymbol{q}_\alpha\, \rho(\boldsymbol{q}_\alpha)\, W(\boldsymbol{k}_\beta,\boldsymbol{q}_\alpha)\, f(\boldsymbol{q}_\alpha) = \nu_\alpha \langle f \rangle_\alpha / \tau_{\beta\alpha}. \tag{4}$$

Here $\alpha,\beta = 1,2$ are band indices; $\langle \cdots \rangle_\alpha$ denotes averaging only over the α band, and $\nu_\alpha = \int d^2\boldsymbol{q}_\alpha\, \rho(\boldsymbol{q}_\alpha) = N_\alpha/N(0)$ are relative densities of states: $N_1 + N_2 = N(0)$, or $\nu_1 + \nu_2 = 1$.

We now assume the order parameters $\Delta(\boldsymbol{k}_\alpha)$ taking constant values Δ_1 and Δ_2 at each of the two bands. Writing Eq. (1) for $\boldsymbol{k}$ in the first band, we have

$$0 = 2\Delta_1 g_1 - 2\omega f_1 + \frac{\nu_1}{\tau_{11}}[g_1\langle f \rangle_1 - f_1 \langle g \rangle_1] + \frac{\nu_2}{\tau_{12}}[g_1\langle f \rangle_2 - f_1\langle g \rangle_2]. \tag{5}$$

For a uniform sample in zero field and with $\boldsymbol{k}$ independent Δ's in each band, the functions f,g are $\boldsymbol{k}$ independent: $\langle f \rangle_\alpha = f_\alpha$ and $\langle g \rangle_\alpha = g_\alpha$. Then, we have

$$0 = \Delta_1 g_1 - \omega f_1 + \nu_2[g_1 f_2 - f_1 g_2]/2\tau_{12}. \tag{6}$$

The equation for the second band differs from this by replacement $1 \leftrightarrow 2$. The fact that τ_{11} and τ_{22} do not enter the system (6) is similar to the case of one-band isotropic material for which nonmagnetic scattering has no effect upon T_c (the Anderson theorem). It is the interband scattering that makes the difference in the two-gap case, the fact stressed already in early work.[21,22]

Two equations (6) are complemented with normalizations $g_\alpha^2+f_\alpha^2=1$ to form a sufficient set. Following Ref. 23, we introduce variables $u_\alpha=g_\alpha/f_\alpha$ and obtain after simple algebra[24,21,22,25]

$$\frac{\omega}{\Delta_1}=u_1+\zeta_1\frac{u_1-u_2}{\sqrt{u_2^2+1}},\qquad \zeta_1=\frac{\nu_2}{2\tau_{12}\Delta_1};$$

$$\frac{\omega}{\Delta_2}=u_2+\zeta_2\frac{u_2-u_1}{\sqrt{u_1^2+1}},\qquad \zeta_2=\frac{\nu_1}{2\tau_{21}\Delta_2}. \tag{7}$$

The Eilenberger functions in terms of variable u are

$$f=1/\sqrt{1+u^2},\qquad g=u/\sqrt{1+u^2}. \tag{8}$$

In general, the system (7) can be solved only numerically. However, near T_c, $u=g/f\gg1$ and one obtains

$$u_1=\frac{\omega}{\Delta_1}\,\frac{\omega+\zeta_1\Delta_1+\zeta_2\Delta_2}{\omega+(\zeta_1+\zeta_2)\Delta_2}; \tag{9}$$

u_2 is obtained by $1\leftrightarrow2$. Clearly, $u_\alpha=\omega/\Delta_\alpha$ in the absence of interband scattering. For $\zeta\gg1$, we have

$$u_1\approx u_2\approx\frac{\omega}{\epsilon^*},\quad \epsilon^*=\frac{(\zeta_1+\zeta_2)\Delta_1\Delta_2}{\zeta_1\Delta_1+\zeta_2\Delta_2}=\langle\Delta\rangle \tag{10}$$

for $\tau_{12}=\tau_{21}$. Moreover, if the interband scattering is strong, Eq. (10) holds at any T. To see this, look for solutions of Eqs. (7) in the form

$$u_\alpha=\frac{\omega}{\epsilon^*}+v_\alpha,\qquad \alpha=1,2, \tag{11}$$

where v_α are small corrections. Substitute these in Eqs. (7) and keep only linear terms in v to obtain

$$v_1=\frac{g^*(\epsilon^*-\Delta_1)}{\Delta_1[1+g^*(\zeta_1+\zeta_2)]}+\frac{\epsilon^*(\zeta_1\Delta_1+\zeta_2\Delta_2)-\Delta_1\Delta_2(\zeta_1+\zeta_2)}{\Delta_1\Delta_2[1+g^*(\zeta_1+\zeta_2)]}, \tag{12}$$

where $g^*=\omega/\sqrt{\omega^2+\epsilon^{*2}}$. For $\zeta_\alpha\to\infty$, v_1 remains small only if ϵ^* is given by expression (10).

It is easy to see that ϵ^* is the common energy gap for both bands' energy gaps. It does not seem possible to provide a general expression for ϵ^* in terms of Δ_α and an arbitrary interband scattering strength. Still, in principle, one can evaluate any thermodynamic property of a two-band material knowing the solutions u of the system (7).

If the ground-state functions (which we call now $f^{(0)}$, $g^{(0)}$) are known, one can study perturbations of the uniform state such as penetration of a weak magnetic field, i.e., the problem of the London penetration depth. The perturbations $f^{(1)}$, $g^{(1)}$ should be found from the full Eilenberger equations;[20] we have for the first band

$$\boldsymbol{v}\boldsymbol{\Pi}f_1=2\Delta_1g_1-2\omega f_1+\frac{\nu_1}{\tau_{11}}[g_1\langle f\rangle_1-f_1\langle g\rangle_1]+\frac{\nu_2}{\tau_{12}}[g_1\langle f\rangle_2-f_1\langle g\rangle_2]. \tag{13}$$

Here, $\boldsymbol{v}$ is the Fermi velocity, $\boldsymbol{\Pi}=\nabla+2\pi i\boldsymbol{A}/\phi_0$. The second equation is obtained by $1\leftrightarrow2$. Two equations for the "anomalous" functions f^+ are obtained from these by complex conjugation and by $\boldsymbol{v}\to-\boldsymbol{v}$.[20] The normalizations $g_\alpha^2+f_\alpha f_\alpha^+=1$ complete the system.

We look for solutions in the form

$$f_\alpha=(f_\alpha^{(0)}+f_\alpha^{(1)})\,e^{i\theta(\boldsymbol{r})},\quad f_\alpha^+=(f_\alpha^{(0)}+f_\alpha^{(1)+})e^{-i\theta(\boldsymbol{r})},$$
$$g_\alpha=g_\alpha^{(0)}+g_\alpha^{(1)},\qquad \alpha=1,2, \tag{14}$$

where $f_\alpha^{(0)}$ and $g_\alpha^{(0)}$ can be expressed in terms of u's obtained solving the system (7). The form (14) takes into account that in the London approximation only the overall phase θ depends on coordinates. We obtain for the corrections after straightforward algebra:

$$g_1^{(1)}\Delta_1'-f_1^{(1)}\omega_1'=if_1^{(0)}\boldsymbol{v}\boldsymbol{P}/2,$$
$$g_1^{(1)}\Delta_1'-f_1^{(1)+}\omega_1'=if_1^{(0)}\boldsymbol{v}\boldsymbol{P}/2,$$
$$2g_1^{(0)}g_1^{(1)}+f_1^{(0)}(f_1^{(1)}+f_1^{(1)+})=0, \tag{15}$$

where $\boldsymbol{P}=\nabla\theta+2\pi\boldsymbol{A}/\phi_0$ and

$$\Delta_1'=\Delta_1+\nu_1f_1^{(0)}/2\tau_{11}+\nu_2f_2^{(0)}/2\tau_{12}, \tag{16}$$

$$\omega_1'=\omega+\nu_1g_1^{(0)}/2\tau_{11}+\nu_2g_2^{(0)}/2\tau_{12}. \tag{17}$$

The equations for the second band (decoupled from the first) are obtained by $1\leftrightarrow2$.

To evaluate the penetration depth we turn to the Eilenberger expression for the current density[20]

$$\boldsymbol{j}=-4\pi|e|N(0)T\,\mathrm{Im}\sum_{\omega>0}\langle\boldsymbol{v}g\rangle, \tag{18}$$

and compare it with the London relation

$$\frac{4\pi}{c}j_i=-(\lambda^2)_{ik}^{-1}\left(\frac{\phi_0}{2\pi}\nabla\theta+\boldsymbol{A}\right)_k. \tag{19}$$

Here, $(\lambda^2)_{ik}^{-1}$ is the tensor of the inverse squared penetration depth; summation over k is implied. We now find $g_1^{(1)}$ from the system (15),

$$g_1^{(1)}=\frac{if_1^{(0)2}\,\boldsymbol{v}\boldsymbol{P}}{2(\Delta_1'f_1^{(0)}+\omega_1'g_1^{(0)})}=i\frac{f_1^{(0)2}g_1^{(0)}}{2\omega_1'}\boldsymbol{v}\boldsymbol{P}; \tag{20}$$

$g_2^{(1)}$ is obtained by replacement $1\leftrightarrow2$. Substituting $g_\alpha^{(1)}$ in Eq. (18) and comparing with Eq. (19) we obtain

$$(\lambda^2)_{ik}^{-1}=\frac{16\pi^2e^2N(0)T}{c^2}\sum_{\alpha,\omega}\nu_\alpha\langle v_iv_k\rangle_\alpha\frac{f_\alpha^2g_\alpha}{\omega_\alpha'}. \tag{21}$$

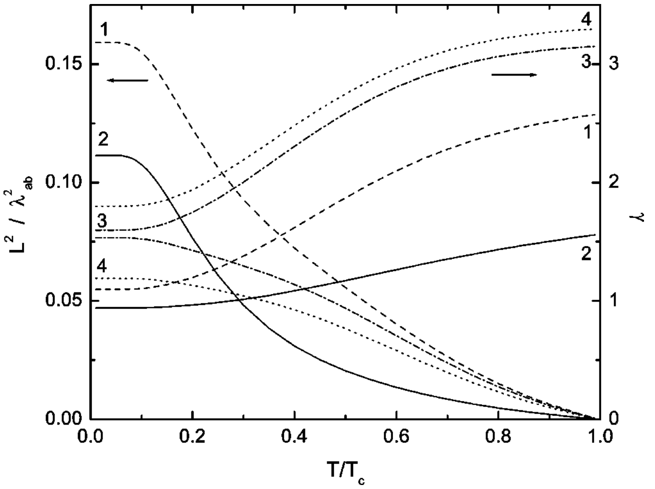

FIG. 1. The anisotropy $\gamma_\lambda = \lambda_c/\lambda_{ab}$ and the inverse square of the penetration depth L^2/λ_{ab}^2 vs T/T_c; $L^2 = 16\pi^2 e^2 N(0)\langle v_a^2\rangle/c^2$. The curves labeled 1 correspond to the clean limit, all $1/\tau$ are zero. The curves labeled 2 and 3 are calculated for a weak interband scattering: $\tau_{12}\Delta_1(0)=500$, $\tau_{21}\Delta_2(0)=2000$ ($\hbar=1$); the curve 2 is for a clean π band, $\tau_{11}\Delta_1(0)=10$, and a dirty σ, $\tau_{22}\Delta_2(0)=0.1$; the curve 3 is for a dirty π and clean σ: $\tau_{11}\Delta_1(0)=0.1$, $\tau_{22}\Delta_2(0)=10$. Curves 4 are for the intermediate interband scattering strength $\tau_{12}\Delta_1(0)=5$, $\tau_{21}\Delta_2(0)=20$, and $\tau_{11}\Delta_1(0)=0.05$, $\tau_{22}\Delta_2(0)=2$.

Only the unperturbed functions f,g enter the penetration depth; for brevity we dropped the superscript (0). Equation (21) is our main result. Thus, to evaluate the penetration depth for given order parameters Δ_α in the presence of scattering, one has to solve the system (7) for $u_\alpha(\omega)$, then to substitute the equilibrium functions f_α, g_α [given in Eq. (8)] in Eq. (21) to sum up over ω.

The band calculations[26] yield for MgB_2: $\nu_1 \approx 0.56$, $\nu_2 \approx 0.44$, $\langle v_a^2\rangle_1 \approx 33.2$, $\langle v_c^2\rangle_1 \approx 42.2$, $\langle v_a^2\rangle_2 \approx 23$, and $\langle v_c^2\rangle_2 \approx 0.5\times 10^{14}$ cm^2/s^2. Tensors $\langle v_i v_k\rangle_1$ and $\langle v_i v_k\rangle_2$ have opposite anisotropies:

$$\frac{\langle v_a^2\rangle_1}{\langle v_c^2\rangle_1} \approx 0.79, \quad \frac{\langle v_a^2\rangle_2}{\langle v_c^2\rangle_2} \approx 46, \tag{22}$$

whereas averaging over the whole Fermi surface yields a nearly isotropic result: $\langle v_a^2\rangle/\langle v_c^2\rangle \approx 1.2$.

In the *clean* limit (all $\tau_{\alpha\beta}\to\infty$) $\omega' = \omega$ and $\Delta'_\alpha = \Delta_\alpha$. Besides, $u_\alpha = \omega/\Delta_\alpha$ and $f_\alpha^2 g_\alpha/\omega'_\alpha = \Delta_\alpha^2/(\omega^2+\Delta_\alpha^2)^{3/2}$. Expression (21) reduces to the result given in Ref. 8. For MgB_2, it gives nearly isotropic penetration depth at low temperatures: at $T=0$ the sums over ω in Eq. (21) are the same; this gives $\gamma_\lambda(0) = \lambda_{cc}/\lambda_{aa} = \sqrt{\langle v_a^2\rangle/\langle v_c^2\rangle} \approx 1.1$. Near T_c, the sums are $\propto \Delta_\alpha^2$, and the contribution of the strongly anisotropic σ band with the large gap dominates; this gives $\gamma_\lambda(T_c) \approx 2.6$. The curve 1 in Fig. 1 shows $\gamma_\lambda(T)$ for this case.

If only the *intraband* scattering is present ($\tau_{12} = \tau_{21} = \infty$), the functions f,g are the same as in the clean limit. We readily obtain

$$\frac{f_\alpha^2 g_\alpha}{\omega'_\alpha} = \frac{\Delta_\alpha^2}{\beta_\alpha^2(\beta_\alpha + 1/2\tau_{\alpha\alpha})} \tag{23}$$

with $\beta_\alpha^2 = \omega^2 + \Delta_\alpha^2$. This expression appears in the standard penetration depth calculations, see, e.g., Ref. 27. For known $\Delta_\alpha(T)$, the sums in Eq. (21) can be evaluated numerically; however, for $T\to 0$, T_c, and in the dirty limit they can be done analytically.

At $T=0$, the sums are replaced with integrals according to $2\pi T\Sigma_\omega = \int_0^\infty d\omega$. Denoting $\mathcal{I}(T) = 2\pi T\Sigma_\omega \Delta^2/\beta^2(\beta + 1/2\tau)$ we obtain $\mathcal{I}(0)=1$ for $\tau\to\infty$ and $\mathcal{I}(0) = \pi\tau\Delta$ for $\tau\Delta \ll 1$.

Near T_c, $g\to 1$ and we have for clean bands $\Sigma_\omega f^2/\omega' = 7\zeta(3)\Delta^2/8\pi^3 T_c^3$, whereas for dirty bands it is $\tau\Delta^2/4T_c^2$.

Different impurities introduced to MgB_2 may affect differently the scattering within the bands.[28–30] It is of interest to see how the anisotropy of λ is affected by differences in scattering times τ_{11} and τ_{22}. We first look at two limiting situations when one of the bands is clean while the other is dirty. If the first (π) band is clean and the second (σ) is a dirty extreme ($\tau_{22}\Delta_2 \to 0$), one can disregard the contribution of the dirty band to obtain for both $T=0$ and $T=T_c$:

$$\gamma_\lambda(0) = \gamma_\lambda(T_c) \approx \sqrt{\frac{\langle v_a^2\rangle_1}{\langle v_c^2\rangle_1}} \approx 0.89. \tag{24}$$

If the π band is dirty and the σ is clean, we have

$$\gamma_\lambda(0) = \gamma_\lambda(T_c) \approx \sqrt{\frac{\langle v_a^2\rangle_2}{\langle v_c^2\rangle_2}} \approx 6.8. \tag{25}$$

These two estimates constitute the minimum and maximum possible values for λ anisotropy of MgB_2. Thus, when one of the bands is clean and the other is dirty we expect a weakly T dependent γ_λ, the value of which is determined by the clean band.

If the intraband scattering is strong in both the bands ($\tau_{11}\Delta_1 \sim \tau_{22}\Delta_2 \ll 1$, $\tau_{12}=\infty$), the bands contribute to the superfluid density tensor $(\lambda^2)_{ik}^{-1}$ as two independent dirty superconductors. To see this, we note that $\omega'_1 \approx g_1/2\tau_{11}$ and the sums over ω in Eq. (21) can be evaluated exactly,

$$\sum_\omega \frac{f_1^2 g_1}{\omega'_1} = \sum_\omega \frac{2\tau_{11}\Delta_1^2}{\omega^2+\Delta_1^2} = \frac{\tau_{11}\Delta_1}{2T}\tanh\frac{\Delta_1}{2T}. \tag{26}$$

Then, we arrive at the result obtained by Gurevich with the help of the dirty limit Usadel equations,[19]

$$(\lambda^2)_{ik}^{-1} = \frac{4\pi^2}{c^2\hbar}\sum_\alpha \sigma_{ik}^{(\alpha)}\Delta_\alpha \tanh\frac{\Delta_\alpha}{2T}, \tag{27}$$

where the anisotropic conductivities of the two bands $\sigma_{ik}^{(\alpha)} = 2e^2\langle v_i v_k\rangle_\alpha \tau_{\alpha\alpha}\nu_\alpha N(0)$ are introduced (we write here $\hbar$ explicitly to avoid confusion in dimensions). This yields,

$$\gamma_\lambda^2(0) = \frac{\sigma_{aa}^{(1)}\Delta_1 + \sigma_{aa}^{(2)}\Delta_2}{\sigma_{cc}^{(1)}\Delta_1 + \sigma_{cc}^{(2)}\Delta_2}, \tag{28}$$

$$\gamma_\lambda^2(T_c) = \frac{\sigma_{aa}^{(1)}\Delta_1^2 + \sigma_{aa}^{(2)}\Delta_2^2}{\sigma_{cc}^{(1)}\Delta_1^2 + \sigma_{cc}^{(2)}\Delta_2^2}. \tag{29}$$

Finally, we discuss the possibility of a *strong interband scattering*. As was shown above, in this case $u=\omega/\epsilon^*+O(1/\zeta)$ in both bands, see Eq. (11). The Eilenberger functions are also the same in the two bands: $f=\epsilon^*/\sqrt{\omega^2+\epsilon^{*2}}$, $g=\omega/\sqrt{\omega^2+\epsilon^{*2}}$, $\epsilon^*=\langle\Delta\rangle$. Evaluation of the sums over ω in Eq. (21) is simple provided the intraband scattering is strong too,

$$(\lambda^2)^{-1}_{ik}=\frac{8\pi^2e^2N(0)}{c^2}\epsilon^*\tanh\frac{\epsilon^*}{2T}\sum_\alpha \nu_\alpha\langle v_iv_k\rangle_\alpha\tau_\alpha, \qquad (30)$$

$$\tau_1=\frac{\tau_{11}\tau_{12}}{\nu_2\tau_{11}+\nu_1\tau_{12}},\quad \tau_2=\frac{\tau_{22}\tau_{21}}{\nu_1\tau_{22}+\nu_2\tau_{21}}. \qquad (31)$$

Thus, all components of $(\lambda^2)_{ik}$ have the same T dependence and the anisotropy parameter is T independent,

$$\gamma_\lambda^2=\frac{\nu_1\langle v_a^2\rangle_1\tau_1+\nu_2\langle v_a^2\rangle_2\tau_2}{\nu_1\langle v_c^2\rangle_1\tau_1+\nu_2\langle v_c^2\rangle_2\tau_2}. \qquad (32)$$

If all τ's are the same, we have $\gamma_\lambda^2=\langle v_a^2\rangle/\langle v_c^2\rangle$. For $T\to T_c$, this result was obtained in Ref. 8; we now see that it holds at any temperature.

To recover the behavior of $\gamma_\lambda(T)$ between 0 and T_c one needs explicit dependencies $\Delta(T)$. Qualitatively, this behavior can be modeled assuming $\Delta_\alpha(T)\approx\Delta_\alpha(0)\tanh\sqrt{1/t-1}$ with, e.g., $\Delta_2(0)=4\Delta_1(0)=2T_c$. Figure 1 shows results of the numerical evaluation of $\gamma_\lambda(T)$ for scattering parameters given in the caption (which are not that extreme as in the above discussion). The curves $\gamma_\lambda(T)$ are obtained by solving Eqs. (7) for u's in two bands and then by evaluation of the sums in Eq. (21). It is worth noting that although the T dependences shown in the figure are obtained using approximate $\Delta(T)$, the end points of these curves at $T=0$ and $T=T_c$ are exact.

We conclude that both the interband and intraband scattering affect strongly the superconducting anisotropy of two-band superconductors in general and of MgB_2 in particular. If one of the MgB_2 bands is dirty, the anisotropy is dominated by a cleaner band: $\gamma_\lambda(T)$ is close to unity (and might be even less than 1) if the π band is in the clean limit, whereas in the opposite situation of a clean σ, $\gamma_\lambda(T)$ is large, being in both cases weakly T dependent. The interband scattering suppresses the T dependence of γ_λ as compared to the clean limit discussed earlier.[8]

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