mate is obtained for the most powerful lasers. The effect is still noticeable for lasers whose power is an order of magnitude lower. The fact that the laser frequency corresponds to the transparency range of a crystal ensures the stability of the effect, at least during short pulses.

If the valence band originates from the f-type states, then not only virtual photoelectrons but virtual photoholes interact strongly with the magnetization of a crystal and, therefore, both types of carrier contribute to the photoferromagnetic effect.

Moreover, the participation of the f electrons in virtual optical transitions should give rise to a dependence of the effect on the nature of polarization of the incident light. In the case of real optical transitions such a dependence is simply a consequence of the spontaneous dichroism of a ferromagnet, caused by the spin–orbit interaction. In the case of virtual transitions, we can speak of virtual dichroism caused by the same interaction. This virtual dichroism is possible also because of the spin–orbit interaction in the conduction band if this band originates from hybridization of s-type atomic states with other states, for example, with d-type states in EuO and EuS.

[1)]The terms with the opposite sign of the frequency are omitted because of the smallness of their contribution to the effect in H_V.

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Differential thermoelectric power of excitonic insulator

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The thermoelectric coefficient η_{ex} of an excitonic insulator was calculated. The case when the electron and hole scattering amplitudes by the impurities are different is considered. It is shown that in this case the temperature dependence of η_{ex} goes through a maximum. It is also shown that allowance for the terms of higher order in the impurity interaction compared with the Born approximation leads to an additional contribution to η_{ex}. This contribution is of zeroth order in T/μ (μ is the chemical potential of the system) and is characteristic of only the exciton phase. As a result, the thermoelectric power of an excitonic insulator is higher than that of semimetals at temperatures close to the transition point.

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Keldysh and Kopaev[1] have shown that semimetals with electron and hole Fermi surfaces of like shape are unstable to formation of electron-hole pairs at arbitrarily weak electron-hole attraction. As a result a transition to the state of an excitonic insulator, with a gap in the energy spectrum, takes place at a certain critical temperature T_d.[2] Allowance for the interband transitions[3] fixes the phase of the order parameter of the excitonic insulator and makes states with homogeneous particle flow impossible. Realignment takes place, but there is no superfluidity. When the electron-hole attraction greatly exceeds the interaction terms responsible for the interband transitions, allowance for the latter leads to small corrections to T_d and to other quantities that describe the realignment.

Although an analogy exists between the excitation spectra of a superconductor and excitonic insulator, the kinetic properties are different in a number of cases. Thus, absorption of ultrasound of frequency $\omega < 2\Delta$ has, in contrast to a superconductor, a maximum below T_d.[4,5] This is due to the fact that the perturbation upsets the electron-hole symmetry, the coherence factors enter with a plus sign and have the same form as for the nuclear-spin relaxation rate in superconductors.

The electric conductivity and thermal conductivity in an excitonic insulator decrease monotonically with decreasing temperature,[6,7] inasmuch as in this case the coherence factors enter with a minus sign. The expression for the thermal conductivity is then analogous to the corresponding expression for superconductors.[8]

Zittartz[7] has shown that in the calculation of the thermoelectric coefficient the coherence factors take the same form as for the thermal conductivity and the electric conductivity. It may seem therefore that the

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thermoelectric coefficient η in the excitonic phase also decreases monotonically below T_d. It is shown in the present paper that there are two causes of the growth of η in an excitonic insulator. The integral with respect to energy for η contains an odd function, and this makes η of the order of T/μ (μ is the chemical potential of the system) when the relaxation time is even in the energy. The value of η is also of the same order in normal metals and in superconductors with nonmagnetic impurities.[9] Deviation from the framework of the Born approximation in the calculation of the scattering amplitude of electrons by magnetic impurities in a normal metal leads to a nonzero differential thermoelectric power already in the zeroth order in T/μ.[10] This is due to the appearance of terms odd in energy in the expression for the relaxation time. The presence of such terms in the scattering amplitude of normal excitations of a superconductor leads to an increase of the thermoelectric coefficient compared with the normal state.[11] The latter is due to the increase of the density of the electronic states near the energy gap, and is not connected with the presence of a homogeneous superconducting current. An analogous phenomenon is therefore possible also in an excitonic insulator and is one of the reasons why η increases below T_d.

It is shown in the present paper that the presence of an energy-odd term in the scattering amplitude causes the thermoelectric power of an excitonic insulator to become of zeroth order in T/μ and to increase at $T<T_d$ when the number of normal excitations is not exponentially small. Such an energy-odd term appears in the scattering amplitudes of normal excitations when account is taken of higher orders in the impurity interaction compared with the Born approximation.[12]

The second cause is that the expression for η acquires additional terms whose coherence factors have a plus sign. These terms are not equal to zero when the electron and hole scattering amplitudes are different. This leads to an increase of η even if the Born approximation is adhered to.

The Hamiltonian of the system is, in the electron-hole representation,[12]

$$H=\sum_{\mathbf{k},\alpha=e,h}\varepsilon_\alpha(\mathbf{k})a_{\alpha\mathbf{k}}^+a_{\alpha\mathbf{k}}+\sum_{\mathbf{k}}(\Delta a_{h\mathbf{k}}^+a_{e-\mathbf{k}}^++\text{H.c})+\sum_{\mathbf{k},\mathbf{k}',\alpha=e,h}J_\alpha(\mathbf{k},\mathbf{k}')a_{\alpha\mathbf{k}}^+a_{\alpha\mathbf{k}'}, \tag{1}$$

where $\varepsilon_\alpha=(k^2-k_{\alpha F}^2)/2m_\alpha$, the term with Δ takes into account the appearance of a "condensate" of electron-hole pairs with zero momentum. The last term of (1) describes intraband scattering by the impurities. It is assumed that this scattering is isotropic. The states described by the Hamiltonian (1) are degenerate with respect to the singlet and triplet structures of the electron-hole pairs, and the spin indices are therefore left out of (1). The number of electrons and holes is assumed equal, so that $k_{eF}=k_{hF}$. The case $k_{eF}\neq k_{hF}$, which is analogous to the presence of a magnetic field in superconductors, was considered in[13] the Born approximation in the impurity interaction for the case equality of the remaining parameters.

We introduce the matrix Green's function

$$\hat{G}_{\mathbf{kk}'}(\omega)=\begin{pmatrix}G_{e\mathbf{kk}'}(\omega) & G_{he-\mathbf{kk}'}(\omega)\\ G^*_{he-\mathbf{kk}'}(\omega) & -G_{h\mathbf{kk}'}(-\omega)\end{pmatrix}, \tag{2}$$

where $G_{\alpha\mathbf{kk}'}(\omega)$ are the Fourier transforms of the normal single-particle retarded Green's functions $G_{\alpha\mathbf{kk}'}(t)=-i\theta(t)\langle[a_{\alpha\mathbf{k}}(t),a^+_{\alpha\mathbf{k}'}(0)]\rangle$, and $G^*_{he-\mathbf{kk}'}(\omega)=-G_{eh-\mathbf{kk}'}(-\omega)$ is the Fourier transform of the anomalous Green's function $G^*_{he-\mathbf{kk}'}(t)=-i\theta(t)\langle[a^+_{h-\mathbf{k}}(t),a^+_{e\mathbf{k}'}(0)]\rangle$.[14] The averaging is over the grand canonical ensemble. Following[12,15] we obtain for the Green's function averaged over the coordinates of the impurity atoms

$$\bar{G}_{\mathbf{kk}'}(\omega)=\delta_{\mathbf{kk}'}\begin{pmatrix}P_1(\tilde{\omega}_h(\omega)+\tilde{\varepsilon}_h(\mathbf{k},\omega)) & -P_1\tilde{\Delta}(\omega)\\ -P_1\tilde{\Delta}^*(\omega) & P_1(\tilde{\omega}_e(\omega)-\tilde{\varepsilon}_e(\mathbf{k},\omega))\end{pmatrix}, \tag{3}$$

where

$$P_1=[(\tilde{\omega}_e-\tilde{\varepsilon}_e)(\tilde{\omega}_h+\tilde{\varepsilon}_h)-|\tilde{\Delta}|^2]^{-1},$$

$$\tilde{\omega}_{e,h}(\omega)=\omega+in\frac{(\tilde{\omega}^2-|\tilde{\Delta}_m|^2)^{1/2}\pi\tilde{\omega}N_{e,h}J^2_{e,h}(1+\pi^2N^2_{h,e}J^2_{h,e})}{(\tilde{\omega}^2-\gamma^2)a}, \tag{4}$$

$$\tilde{\Delta}^*(\omega)=\Delta^*+\frac{2nN_e\pi\tilde{\Delta}^*J_eJ_hm_h}{(m_e+m_h)(\tilde{\omega}^2-\gamma^2)a}[\pi\tilde{\omega}(J_eN_e-J_hN_h)+i(\tilde{\omega}^2-|\tilde{\Delta}_m|^2)^{1/2}(1+\pi^2N_eN_hJ_eJ_h)], \tag{5}$$

$$\chi_{e,h}(\omega)=\frac{\pm n\pi^2\tilde{\omega}^2J_eJ_hN_{h,e}(J_eN_e-J_hN_h)}{(\tilde{\omega}^2-\gamma^2)a}-\frac{nJ_{e,h}(\gamma^2-|\tilde{\Delta}_m|^2)(1+\pi^2N_eN_hJ_eJ_h)}{(\tilde{\omega}^2-\gamma^2)a}, \tag{6}$$

$$a=(1+\pi^2N_eN_hJ_eJ_h)^2+\pi^2(J_eN_e-J_hN_h)^2, \tag{7}$$

$$\gamma^2=|\tilde{\Delta}_m|^2(1+\pi^2N_eN_hJ_eJ_h)^2/a, \tag{8}$$

$$\tilde{\varepsilon}_{e,h}(k,\omega)=\varepsilon_{e,h}(k)-\chi_{e,h}(\omega), \tag{9}$$

$$\tilde{\omega}=\frac{(\tilde{\omega}_e+\chi_e)m_e+(\tilde{\omega}_h-\chi_h)m_h}{m_e+m_h}, \tag{10}$$

$$\tilde{\Delta}_m=2\tilde{\Delta}(m_em_h)^{1/2}/(m_e+m_h),\quad \Delta_m=2\Delta(m_em_h)^{1/2}/(m_e+m_h), \tag{11}$$

$N_{e,h}$ is the state density of the electrons or holes of the semimetal on the Fermi surface, and n is the concentration of the impurity atoms.

Gradients of the temperature and of the electric field produce in an excitonic insulator a normal-excitation current

$$\mathbf{j}=\sigma_{ex}\mathbf{E}-\eta_{ex}\nabla T. \tag{12}$$

In an open conductor $j=0$, and the thermoelectric field is then given by

$$\mathbf{E}=(\eta_{ex}/\sigma_{ex})\nabla T. \tag{13}$$

The conductivity σ_{ex} of an excitonic insulator was calculated in[6]. It was shown that at $T<T_d$ the conductivity σ_{ex} is lower than the semimetal conductivity σ_{sm}.

To calculate the coefficient η_{ex} we follow a procedure similar to the one used to obtain η for a superconductor.[11] The difference is that in this case there are two types of carrier, and we obtain

$$\eta_{ex}=-\frac{\hbar^3 e}{24T^2k_B}\left\{\frac{2m_e}{m_e+m_h}\int_{-\infty}^{\infty}\frac{d\omega\,\omega v_e^2(\mu_e+\omega)N_e(\mu_e+\omega)}{\mathrm{ch}^2(\beta\omega/2)\,\mathrm{Im}\,\varepsilon_{e1}(\omega)\,\mathrm{Im}\,\varepsilon_{e2}(\omega)}\right.$$

$$\times\left[\mathrm{Im}\,(\tilde\omega^2-|\tilde\Delta_m|^2)^{1/2}\left(1+\frac{|\tilde\omega^2|-|\tilde\Delta_m|^2}{|\tilde\omega^2-|\tilde\Delta_m|^2|}\right)-\mathrm{Im}\,d_1(\omega)\cdot 2\,\mathrm{Re}\frac{\tilde\omega}{(\tilde\omega^2-|\tilde\Delta_m|^2)^{1/2}}\right]$$

$$-\frac{2m_h}{m_e+m_h}\int_{-\infty}^{\infty}\frac{d\omega\,\omega v_h^2(\mu_h+\omega)N_h(\mu_h+\omega)}{\mathrm{ch}^2(\beta\omega/2)\,\mathrm{Im}\,\varepsilon_{h1}(\omega)\,\mathrm{Im}\,\varepsilon_{h2}(\omega)}$$

$$\left.\times\left[\mathrm{Im}\,(\tilde\Omega^2-|\tilde\Delta_m|^2)^{1/2}\left(1+\frac{|\tilde\Omega|^2-|\tilde\Delta_m|^2}{|\tilde\Omega^2-|\tilde\Delta_m|^2|}\right)-\mathrm{Im}\,d^2\cdot 2\,\mathrm{Re}\frac{\tilde\Omega}{(\tilde\Omega^2-|\tilde\Delta_m|^2)^{1/2}}\right]\right\}$$

where

$$\mu_{e,h}=k_F^2/2m_{e,h}, \tag{14}$$

$$\tilde\Omega=[m_e(\tilde\omega_e-\chi_e)+m_h(\tilde\omega_h+\chi_h)]/(m_e+m_h), \tag{15}$$

$$d_1=[m_e(\tilde\omega_e+\chi_e)-m_h(\tilde\omega_h-\chi_h)]/(m_e+m_h), \tag{16}$$

$$d_2=[m_h(\tilde\omega_h+\chi_h)-m_e(\tilde\omega_e-\chi_e)]/(m_e+m_h), \tag{16'}$$

$$\varepsilon_{e1,2}=d_1\pm(\tilde\omega^2-|\tilde\Delta_m|^2)^{1/2},\quad \varepsilon_{h1,2}=d_2\pm(\tilde\Omega^2-|\tilde\Delta_m|^2)^{1/2}. \tag{17}$$

In the case of a semimetal, when $\Delta=0$ and $\mathrm{Im}\chi_{e,h}=0$, we obtain in the model of free electrons and holes the known expression

$$\eta_{sm}=\frac{2}{9}\hbar^3eTk_B^2\pi^2\left[\frac{\partial}{\partial\omega}(N_e(\mu_e+\omega)v_e^2(\mu_e+\omega)\tau_e(\mu_e+\omega))|_{\omega=0}\right.$$

$$\left.-\frac{\partial}{\partial\omega}(N_h(\mu_h+\omega)v_h^2(\mu_h+\omega)\tau_h(\mu_h+\omega))|_{\omega=0}\right], \tag{18}$$

where $\mathrm{Im}\tilde\omega_{e,h}=\frac{1}{2}\tau_{e,h}=n\pi J_{e,h}^2N_{e,h}$

Expression (14) differs from those obtained in [7,13] in that it contains terms proportional to $\mathrm{Im}d_1$ and $\mathrm{Im}d_2$, whose coherence factors have a plus sign. It is these terms which cause the thermoelectric power to increase in the Born approximation, when $\chi_{e,h}=0$, $\gamma=\tilde\Delta_m$, $\tilde\Omega=\tilde\omega=\tilde z=(\tilde\omega_e m_e+\tilde\omega_h m_h)/(m_e+m_h)$, and $\mathrm{Im}d_{1,2}=\mathrm{Im}[(m_{e,h}\tilde\omega_{e,h}-m_{h,e}\tilde\omega_{h,e})/(m_e+m_h)]$. The last quantities are not equal to zero when $J_{e,h}N_{e,h}m_{e,h}\neq J_{h,e}N_{h,e}m_{h,e}$. It is seen from (14) that when all the parameters of the e and h bands are equal the quantity η_{ex} vanishes. Therefore at nonzero η_{ex} there are always band parameters that are not equal and the result is $\mathrm{Im}d_{1,2}\neq 0$. For the sake of argument we assume that $m_e\neq m_h$.

The integral with respect to ω in (14) is calculated by the method developed in [16] for the investigation of the densities of the electronic states of superconducting alloys. We assume that

$$\frac{N_em_e-N_hm_h}{2N_{e,h}m_{e,h}}(n(\omega)-1)\ll 1,\quad n(\omega)=\mathrm{Re}\frac{u(\omega)}{[u^2(\omega)-1]^{1/2}}. \tag{19}$$

The equations for $u(\omega)$ are similar to those obtained in [17] for superconductors with magnetic impurities:

$$u=\omega/\Delta_m+i\sigma u/(u^2-1)^{1/2},\quad u=\tilde z(\omega)/\tilde\Delta_m(\omega).$$

The pair-breaking parameter σ for a neutral impurity ($J=J_e=J_h$) is of the form

$$\sigma=\frac{n\pi J^2(m_e-m_h)(N_e-N_h)}{\Delta_m(m_e+m_h)}.$$

For a charged impurity ($j_e=-J_h=J$), which can be considered when the electron and hole numbers are equal only in the presence of a compensating impurities, we obtain for [6]

$$\sigma=n\pi J^2(N_e+N_h)/\Delta_m.$$

Taking into account the inequality (19), we get from (14)

$$\eta_{exB}\cong\eta_{sm}+\frac{\hbar^3e}{3T^2k_B}A\int_0^{\infty}\frac{d\omega\,\omega^2}{\mathrm{ch}^2(\beta\omega/2)}(n^2(\omega)-1), \tag{20}$$

where

$$A=\frac{(N_hm_h-N_em_e)}{2m_hN_h}\frac{\partial}{\partial\omega}[N_e(\mu_e+\omega)v_e^2(\mu_e+\omega)\tau_e(\mu_e+\omega)]|_{\omega=0}$$

$$+\frac{N_hm_h-N_em_e}{2m_eN_e}\frac{\partial}{\partial\omega}[N_h(\mu_h+\omega)v_h^2(\mu_h+\omega)\tau_h(\mu_h+\omega)]|_{\omega=0}. \tag{21}$$

We break up the integration region in (20) into two intervals and substitute in each of them the corresponding values for $n(\omega)$.[16] As a result we obtain at $\Delta_m\beta\sigma^{2/3}\ll 1$ and $\sigma\ll 1$.

$$\eta_{exB}=\eta_{sm}+\frac{\hbar^3e}{6}Tk_B^2A(\Delta_m\beta)^3\frac{\ln(4/3\sigma^{3/2})}{\mathrm{ch}^2(\beta\Delta_m/2)}. \tag{22}$$

The quantity η_{exB} does not depend on the sign of the interaction J. In the free-electron model the sign of N_{sm} is determined by the factor $(1/m_e-1/m_h)$ and coincides with the sign of the coefficient A. In other words, the thermoelectric coefficient is larger in an electronic insulator than in a semimetal. Let us estimate the second term in (22). At $m_h-m_e\sim m_e/100$, $\Delta_m\beta\sim 1$, and $\sigma\sim 10^{-5}$ we find that it is of the order of η_{sm} and increases with increasing concentration η relative to η_{sm}. But n is bounded from below by the inequality (19). As $n\to 0$ formulas (20) and (22) do not hold, and the term supplementing η_{sm} is of the order of $\sigma^{2/3}\ll 1$, a natural result.

Expression (22) has no analog in the case of a superconductor, since the amplitude of quasiparticle scattering by a magnetic impurity in a superconductor, unlike in an excitonic insulator, is the same because of the averaging over the spins of the atoms.[15,17] This, together with the growth of the densities of the electronic states near the energy gap, leads to $\eta_{exB}>\eta_{sm}$.

Allowance for terms of higher order in the impurity interaction compared with the Born approximation leads to an additional contribution to the coefficient η_{ex}. This contribution, $\eta_{ex\,ad}$, just as in the case of a superconductor,[11] differs from zero in the zeroth order in T/μ and is due to the fact that $\mathrm{Im}\chi_{e,h}(\omega)=-\mathrm{Im}\chi_{e,h}(-\omega)$ are not equal to zero with $m_e\neq m_h$.

Removing in (14) all the smooth functions N_e,N_h,v_e,v_h to the Fermi surface and using the properties $\tilde\omega(\omega)=-\tilde\Omega^*(-\omega)$, $d_1(\omega)=d_2^*(-\omega)$, $\tilde\Delta_m^*(\omega)=\tilde\Delta_m^*(-\omega)$, and also assuming that $N_ev_e^2m_e=N_hv_h^2m_h$, we get

$$\eta_{ex\,ad}=-\frac{\hbar^3eN_ev_e^2}{6T^2k_B}\frac{m_e}{m_e+m_h}\int_0^{\infty}\frac{d\omega\,\omega\,\mathrm{Im}\,(\tilde\omega^2-|\tilde\Delta_m|^2)^{1/2}}{\mathrm{ch}^2(\beta\omega/2)\,\mathrm{Im}\,\varepsilon_{e1}(\omega)\,\mathrm{Im}\,\varepsilon_{e2}(\omega)}$$

$$\times\left(1+\frac{|\tilde\omega(\omega)|^2-|\tilde\Delta_m|^2}{|\tilde\omega^2(\omega)-|\tilde\Delta_m|^2|}\right)+\frac{\hbar^3eN_hv_h^2m_h}{6T^2k_B(m_e+m_h)}$$

$$\times\int_0^{\infty}\frac{d\omega\,\omega\,\mathrm{Im}\,(\tilde\Omega^2-|\tilde\Delta_m|^2)^{1/2}}{\mathrm{ch}^2(\beta\omega/2)\,\mathrm{Im}\,\varepsilon_{h1}(\omega)\,\mathrm{Im}\,\varepsilon_{h2}(\omega)}\left(1+\frac{|\tilde\Omega(\omega)|^2-|\tilde\Delta_m|^2}{|\tilde\Omega^2(\omega)-|\tilde\Delta_m|^2|}\right) \tag{23}$$

It is seen from the last formula that at $\chi_{e,h}=0$ we have $\eta_{ex\,ad}=0$.

To estimate the integral with respect to ω in (23), we

assume that the deviation from the Born approximation is small. Then formula (6) for $\chi_{e,h}$ becomes

$$\chi_{e,h}(\omega)=\frac{\pm n\pi^2 J^3 N_{h,e}(N_e-N_h)-nJ(\gamma^2-|\tilde{\Delta}_m|^2)}{\tilde{z}^2-|\tilde{\Delta}_m|^2}$$

Expanding all the quantities in (23) in powers of $\chi_{e,h}$ and using the inequality (19), we get

$$\eta_{ex\,ad}=\frac{4}{3}\frac{\hbar^3 e N_e v_e^2 m_e}{T^2 k_B (m_e+m_h)^2}\int_0^\infty \frac{\omega d\omega}{\mathrm{ch}^2(\beta\omega/2)\,\mathrm{Im}\,(\tilde{z}^2-|\tilde{\Delta}_m|^2)^{1/2}} \tag{24}$$
$$\times\frac{|\tilde{\Delta}_m|^2\,\mathrm{Im}\,\tilde{z}}{|\tilde{z}^2-|\tilde{\Delta}_m|^2|}\,\mathrm{Im}\frac{\chi_e m_e-\chi_h m_h}{\tilde{z}^2-|\tilde{\Delta}_m|^2}.$$

Calculation of the integral with respect to ω in (24) is by the method developed in [16]. The result is

$$\eta_{ex}=\eta_{exB}+\eta_{ex\,ad} \tag{25}$$

where

$$\eta=\frac{\hbar^3 e k_B N_e^2 v_e^2 m_e m_h \pi^2 n J^3 (N_h-N_e)\tau^2(\Delta_m\beta)^2}{(m_e+m_h)^2\sigma\,\mathrm{ch}^2(\beta\Delta_m/2)}, \tag{26}$$
$$\frac{1}{2\tau}=\pi n J^2\frac{(N_e m_e+N_h m_h)}{m_e+m_h}.$$

The sign of $\eta_{ex\,ad}$ at $J_e=J_h=J>0$ is the same as the sign of η_{exB}. Thus, in the presence of a neutral impurities the expression for the thermoelectric coefficient η_{ex} acquires an additional term and the thermoelectric power increases. The ratio of $\eta_{ex\,ad}$ to η_{sm} is of the order of

$$\frac{\mu_{e,h}}{Tk_B}(\Delta_m\beta)^2\frac{N_{e,h}J_{e,h}}{\sigma}$$

and can be much larger than unity. In the case of a charged impurity the sign of $\eta_{ex\,ad}$ is determined by the sign of J, so that in the case considered by us, when a compensating impurity is present, $\eta_{ex\,ad}=0$. This is analogous to normal metals[11]: anomalous values of the thermoelectric power appear only if simultaneous account is taken of both the exchange and the potential interactions of the electrons with the impurity.

In the calculation of the thermal conductivity and electric conductivity, terms of the type $\mathrm{Im}\,d_1$ and $\mathrm{Im}(\chi_e-\chi_h)$ cancel out. All that are left are the terms proportional to $\mathrm{Im}(\chi_e+\chi_h)$, which give increments of the order of T/μ, so that the thermal conductivity and electric conductivity decrease monotonically below T_d.[6,7]

Formulas (22) and (26) were obtained for the temperature region defined by the inequality $\Delta_m\beta\sigma^{2/3}\ll 1$. At $T\ll T_d$, when the number of excitations is exponentially small, the formulas are not valid and in this case η_{ex} is also exponentially small. Substituting (22) and (26) in (13) we find that the thermoelectric field for an excitonic insulator is larger than that of a semimetal in the region where $\Delta_m\beta\sigma^{2/3}\ll 1$.

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