Spin-galvanic effect due to optical spin orientation

S.D. Ganichev\textsuperscript{1,2}, Petra Schneider\textsuperscript{1}, V.V. Bel’kov\textsuperscript{2}, E.L. Ivchenko\textsuperscript{2}, S.A. Tarasenko\textsuperscript{2}, W. Wegscheider\textsuperscript{1}, D. Weiss\textsuperscript{1}, D. Schuh\textsuperscript{3}, B.N. Murdin\textsuperscript{4}, P.J. Phillips\textsuperscript{5}, C.R. Pidgeon\textsuperscript{5}, D.G. Clarke\textsuperscript{4}, M. Merrick\textsuperscript{4}, P. Murzyn\textsuperscript{5}, E.V. Beregulin\textsuperscript{2}, and W. Prettl\textsuperscript{1}

\textsuperscript{1} Fakultät Physik, University of Regensburg, 93040, Regensburg, Germany
\textsuperscript{2} A.F. Ioffe Physico-Technical Institute, 194021 St. Petersburg, Russia
\textsuperscript{3} Walter Schottky Institute, TU Munich, D-85748 Garching, Germany
\textsuperscript{4} University of Surrey, Guildford, GU2 7XH, UK
\textsuperscript{5} Department of Physics, Heriot-Watt University, Edinburgh, UK

(Dated: November 17, 2018)

Abstract

Under oblique incidence of circularly polarized infrared radiation the spin-galvanic effect has been unambiguously observed in (001)-grown \textit{n}-type GaAs quantum well (QW) structures in the absence of any external magnetic field. Resonant inter-subband transitions have been obtained making use of the tunability of the free-electron laser FELIX. It is shown that a helicity dependent photocurrent along one of the \textlangle{110}\textrangle axes is predominantly contributed by the spin-galvanic effect while that along the perpendicular in-plane axis is mainly due to the circular photogalvanic effect. This strong non-equivalence of the [110] and [110] directions is determined by the interplay between bulk and structural inversion asymmetries. A microscopic theory of the spin-galvanic effect for direct inter-subband optical transitions has been developed being in good agreement with experimental findings.
The spin of electrons and holes in solid state systems is an intensively studied quantum mechanical property showing a large variety of interesting physical phenomena. Lately there is much interest in the use of the spin of carriers in semiconductor heterostructures together with their charge for novel applications like spintronics [1]. The necessary conditions to realize spintronic devices are high spin polarizations in low dimensional structures and large spin-splitting of subbands in \( k \)-space. The latter is important for the ability to control spins with an external electric field by the Rashba effect [2]. Significant progress has been achieved recently in generating large spin polarizations, in demonstrating the Rashba splitting and also in using the splitting for manipulating the spins [1]. At the same time as these conditions are fulfilled it has been shown that the spin polarization itself drives a current if the spins are oriented in the plane of the quantum well (QW) [3]. This spin-galvanic effect was previously demonstrated with optical excitation and the assistance of an external magnetic field to achieve an in-plane polarization. As a step towards the long term aim of showing its existence with only electric injection we report here the demonstration of the optically induced spin-galvanic effect in zero magnetic field. We also present the microscopic theory of this effect.

The spin-galvanic effect has been observed at room temperature by studying transitions between size quantized subbands \( e_1 \) and \( e_2 \) in \( n \)-type GaAs/AlGaAs quantum wells (QW). Typical samples, grown along \( z \parallel [001] \) by molecular-beam-epitaxy, consisting of 30 QWs with a well width of 7.6 nm, 8.2 nm and 8.6 nm, and free-carrier density in a single well \( n_e \) of about \( 2 \cdot 10^{12} \) cm\(^{-2} \) were investigated at room temperature. Samples were quadratic in shape, with edges oriented along the \( x \parallel [1\bar{1}0] \) and \( y \parallel [110] \) crystallographic directions. Two pairs of ohmic contacts were attached in the center of opposite sample edges (see Fig. 1).

In Fig. 2 the absorption spectrum of the sample containing 8.2 nm wide QWs obtained by Fourier transform spectroscopy in a waveguide geometry is shown by the dotted curve. The \( e_1 \) to \( e_2 \) resonance occurs at the photon energy 130.4 meV and the full width half maximum was 16.8 meV. In order to excite resonantly and to obtain a measurable photocurrent it was necessary to have a tunable high power radiation
source for which we used the free electron laser "FELIX" at FOM-Rijnhuizen in The Netherlands [4]. The output pulses of light from FELIX were chosen to be 3 ps long, separated by 40 ns, in a train (or "macropulse") of duration of 5 µs. The macropulses had a repetition rate of 5 Hz.

On illumination of the QW structures by circularly polarized radiation at oblique incidence in (xz)- or (yz)-plane a current signal perpendicular to the plane of incidence was measured, e.g. in y direction for the configuration depicted in Fig. 1. Left handed (\(\sigma^-\)) and right handed (\(\sigma^+\)) circularly polarized radiation was achieved making use of a Fresnel rhomb. The photocurrent signals generated in the unbiased devices at room temperature were measured via an amplifier with a response time of the order of 1 µs, i.e. averaged over the macropulse. The voltage in response to a laser pulse was recorded by an oscilloscope.

The observed current is proportional to the helicity \(P_{\text{circ}}\) of the radiation. The photon energy dependence of the current was measured for incidence in two different planes with in-plane component of propagation along the x and y directions. In Fig. 2, the observed current for both directions is plotted as a function of photon energy \(\hbar \omega\) for \(\sigma^+\) polarized radiation together with the absorption spectrum. It can be seen that for current along \(x \parallel [1\bar{1}0]\) the shape is similar to the derivative of the absorption spectrum, and in particular there is a change of sign which occurs at the line center of
the absorption. When the sample was rotated by 90° about \( z \), so that light propagates now along \( x \) and the current flows along \( y \parallel [110] \), the sign reversal in the current disappears and its shape follows more closely the absorption spectrum.

FIG. 2: Photocurrent in QWs normalized by the light power \( P \) at oblique incidence of right-handed circularly polarized radiation on \( n \)-type (001)-grown GaAs/AlGaAs QWs of 8.2 nm width at \( T = 293 \) K as a function of the photon energy \( \hbar \omega \). Circles: current in [110] direction in response to irradiation parallel [1¯10]. Squares: current in [110] direction in response to irradiation parallel [110]. The dotted line shows the absorption measured using a Fourier transform spectrometer.

It has been shown in [3, 5] that in quantum wells belonging to one of the gyrotropic crystal classes a non-equilibrium spin polarization of electrons uniformly distributed in space causes a directed motion of electrons in the plane of the QW. On a microscopic level spin photocurrents are the result of spin orientation in systems with \( k \)-linear terms in the electron effective Hamiltonian which are characteristic of gyrotropic me-
In general, two mechanisms contribute to spin photocurrents: photoexcitation and scattering of photoexcited carriers. The first effect is the spin orientation induced circular photogalvanic effect (CPGE) which is caused by an asymmetry of the momentum distribution of carriers excited in optical transitions \[5, 6\]. The second effect is the spin-galvanic effect which in general does not need optical excitation but is a result of an asymmetric spin relaxation \[3\]. The current due to CPGE is spin polarized and decays with the momentum relaxation time \(\tau_p\) of photoexcited free carriers whereas the spin-galvanic effect induced current is unpolarized but decays with the spin relaxation time \(\tau_s\). Both effects are illustrated in Fig. 3.

The change of sign of the photocurrent with photon energy is characteristic for CPGE at resonant transitions in \(n\)-type QWs and has been described previously \[6\]. As illustrated in Fig. 3a for \(\sigma_+\) radiation and at a small photon energy less than the energy separation between \(e_1\) and \(e_2\) at \(k_x = 0\), excitations occur preferentially at positive \(k_x\). We note that for \(C_{2v}\) symmetry the optical transitions are spin-conserving but spin-dependent \[6\]. This causes a stronger reduction in the electron population at positive \(k_x\) in the lower \(|-1/2\rangle_y\)-subband and therefore a spin-polarized current in positive \(x\) direction. We note that there is a corresponding increase of the electron population in the \(e_2\) \(|-1/2\rangle_y\)-subband, also asymmetrical, but in our case this randomizes quickly via optical phonon scattering and therefore does not contribute significantly to the current \[6\]. Increase of the photon energy shifts the dominating transition towards negative \(k_x\) and reverses the current. In fact it has been shown that the CPGE at intersubband absorption in \(n\)-type QWs is proportional to the derivative of the absorption spectrum \[6\]. This behaviour is observed for the current in \(x \parallel [\bar{1}10]\) direction. In particular, the position of the sign inversion of the current coincides with the maximum of the absorption spectrum which shows that the spin-galvanic effect for this direction is vanishingly small and the current is caused by the CPGE.

In contrast to the CPGE the sign of the spin-galvanic current is independent of the wavelength \[7\]. This can be seen from Fig. 3b which illustrates the origin of the spin-galvanic effect. All that is required is a spin orientation of the lower subband, and asymmetrical spin relaxation then drives a current \[3\]. In our case the spin orientation
FIG. 3: Microscopic picture of (a) circular photogalvanic effect and (b) spin-galvanic effect at inter-subband excitation in C$_{2v}$ point group samples. In (a) the current $j_x$ is caused by the imbalance of optical transition probabilities at $k_x^-$ and $k_x^+$ decaying with the momentum relaxation time $\tau_p$. Excitation with $\sigma_+$ radiation of $\hbar \omega$ less than the energy subband separation at $k=0$, $\varepsilon_{21}$, induces direct spin-conserving transitions (vertical arrows) at $k_x^-$ and $k_x^+$. The rates of these transitions are different as illustrated by the different thicknesses of the arrows (reversing the angle of incidence mirrors the transition rates). This leads to a photocurrent due to an asymmetrical distribution of carriers in $k$-space if the splitting of the $e_1$ and $e_2$ subbands is non-equal. Increase of the photon energy shifts more intensive transitions to the left and less intensive to the right resulting in a current sign change. In (b) the current occurs after thermalization in the lowest subband which results in the spin orientation in the $e_1$ subband. This spin-galvanic current is caused by asymmetric spin-flip scattering. The rate of spin-flip scattering depends on the value of the initial and final $k$-vectors. Thus transitions sketched by dashed arrows yield an asymmetric occupation of both subbands and hence a current flow which decays with the spin relaxation time $\tau_s$. The magnitude of the spin polarization and hence the current depends on the initial absorption strength but not on the momentum $k$ of the transition. Therefore the shape of the spectrum of the spin-galvanic current follows the absorption.
is generated by resonant spin-selective optical excitation followed by spin-non-specific thermalization. The magnitude of the spin polarization and hence the current depends on the initial absorption strength but not on the momentum $k$ of transition. Therefore there is no sign change and the shape of the spectrum follows the absorption. The lack of a sign change for current along $y \parallel [110]$ in the experiment shows that the spin-galvanic dominates for this orientation.

In order to understand the difference between the two orientations we now introduce a phenomenological picture for the $C_{2v}$ symmetry representing samples investigated here. Phenomenologically the spin-galvanic effect (SGE) and the circular photogalvanic effect in $x$ and $y$ directions are given by

$$j_{SGE,x} = Q_{xy} S_y, \quad j_{SGE,y} = Q_{yx} S_x.$$  (1)

$$j_{CPGE,x} = \gamma_{xy} \hat{e}_y E_0^2 P_{circ}, \quad j_{CPGE,y} = \gamma_{yx} \hat{e}_x E_0^2 P_{circ}.$$  (2)

where $j$ is the photocurrent density, $Q$ and $\gamma$ are second rank pseudo-tensors, $S$ is the average spin of electrons in QWs, $E_0$, $P_{circ}$ and $\hat{e}$ are the amplitude of the electromagnetic wave, the degree of circular polarization and the unit vector pointing in the direction of light propagation, respectively. In the present case $S$ is obtained by optical orientation, its sign and magnitude are proportional to $P_{circ}$ and it is oriented along the in-plane component of $\hat{e}$ (see Fig. 1). Because of tensor equivalence of $Q$ and $\gamma$ the spin-galvanic current induced by circularly polarized light always occurs simultaneously with the CPGE. If the in-plane component of $\hat{e}$ is oriented along $[1 \bar{1} 0]$ or $[110]$, i.e. $x$ or $y$, then both currents flow normal to the light propagation direction. The strength of the current is different for the radiation propagating along $x$ or $y$. This is due to the non-equivalence of the crystallographic axes $[1 \bar{1} 0]$ and $[110]$ because of the two-fold rotation axis in $C_{2v}$ symmetry.

Both currents are caused by spin splitting of subbands in the $k$-space. This splitting is due to $k$-linear terms in the Hamiltonian of the form $\hat{H}' = \sum_{lm} \beta_{lm} \sigma_l k_m$, where $\beta_{lm}$ is a second rank pseudo-tensor and $\sigma_l$ are the Pauli-matrices. The tensors $\gamma$ and $Q$ determining the current are related to the transposed pseudo-tensor $\beta$. They are subjected to the same symmetry restrictions so that their irreducible components differ.
FIG. 4: Schematic 2D band structure with $k$-linear terms for $C_{2v}$ symmetry. The energy $\varepsilon$ is plotted as a function of $k_x$ and $k_y$ for equal strength of the BIA and SIA terms in the Hamiltonian. The bottom plate shows the distribution of spin orientations at the 2D Fermi energy: (b) for BIA and SIA terms with equal strength and (c) for different strength of SIA and BIA terms in the Hamiltonian. The difference of subband spin-splitting in $x$ and $y$ directions, which is clearly seen from (b) and (c) sketches, occurs due to non-equality of $\beta_{xy}\sigma_x k_y$ and $\beta_{yx}\sigma_y k_x$. These component of $\beta$ may also be written as $\beta_{xy} = (\beta_{BIA} + \beta_{SIA})/2$ and $\beta_{yx} = (\beta_{BIA} - \beta_{SIA})/2$, respectively. Arrows indicate the orientation of spins. Only by scalar factors. The non-zero components of the pseudo-tensor $\beta_{lm}$ depend on the symmetry and the coordinate system used. For (001)-crystallographic orientation grown QWs of $C_{2v}$ symmetry and in the coordinate system $(xyz)$, relevant to our experimental set-up, there are two non-zero tensor elements $\beta_{xy} \neq \beta_{yx}$ which may also be different for $e_1$- and $e_2$ subbands. It is reasonable to introduce symmetric and anti-symmetric tensor components $\beta^{(\nu)}_{BIA} = (\beta_{xy}^{(\nu)} + \beta_{yx}^{(\nu)})/2$ and $\beta^{(\nu)}_{SIA} = (\beta_{xy}^{(\nu)} - \beta_{yx}^{(\nu)})/2$, where $\nu=1,2$ indicates the $e_1$ and $e_2$ subbands respectively. Here $\beta^{(\nu)}_{BIA}$ and $\beta^{(\nu)}_{SIA}$ result from bulk inversion asymmetry (BIA) also called the Dresselhaus term [8] (including a possible interface inversion asymmetry [9]) and from structural inversion asymmetry (SIA) usually called the Rashba term [2], respectively. In order to illustrate band structures with a $k$-linear term in Fig. 4 we plotted the energy $\varepsilon$ as a function of $k_x$ and $k_y$ and constant energy surfaces for different relations between $\beta_{BIA}$ and $\beta_{SIA}$ which are assumed to be positive. The non-equivalence of $x$ and $y$ directions for $|\beta_{xy}| \neq |\beta_{yx}|$ is clearly seen from Fig. 4.
As discussed above and sketched in Fig. 3 both CPGE and spin-galvanic currents, say in $x$ direction, are caused by the band splitting in $k_x$ direction and therefore are proportional to $\beta_{yx}$ (for current in $y$-direction one should interchange the indices $x$ and $y$). Then the currents in the $x$ and $y$ directions read

$$j_x = A_{\text{CPGE}}[(\beta_{\text{BIA}}^{(1)} - \beta_{\text{SIA}}^{(1)}) - (\beta_{\text{BIA}}^{(2)} - \beta_{\text{SIA}}^{(2)})]P_{\text{circ}}\hat{e}_y + A_{\text{SGE}}(\beta_{\text{BIA}}^{(1)} - \beta_{\text{SIA}}^{(1)})S_y$$

(3)

and

$$j_y = A_{\text{CPGE}}[(\beta_{\text{BIA}}^{(1)} + \beta_{\text{SIA}}^{(1)}) - (\beta_{\text{BIA}}^{(2)} + \beta_{\text{SIA}}^{(2)})]P_{\text{circ}}\hat{e}_x + A_{\text{SGE}}(\beta_{\text{BIA}}^{(1)} + \beta_{\text{SIA}}^{(1)})S_x$$

(4)

where $A_{\text{CPGE}}$ and $A_{\text{SGE}}$ are factors related to $\gamma$ and $Q$, respectively. The magnitude of the CPGE is determined by the value of $k$ in the initial and final states, and hence on the spin splitting ($\beta_{\text{BIA}}$ and $\beta_{\text{SIA}}$) of both $e_1$ and $e_2$ subbands. In contrast, the spin-galvanic effect is due to relaxation between the spin states of the lowest subband $e_1$ and hence only on $\beta_{\text{BIA}}^{(1)}$ and $\beta_{\text{SIA}}^{(1)}$. The Eqs. (3) and (4) show that in directions $x$ and $y$ the spin-galvanic effect and the CPGE are proportional to terms containing the difference and the sum, respectively, of BIA and SIA terms. When they add (see Eq. 4) it appears in our samples that the spin-galvanic effect dominates over the CPGE which is proved by the lack of sign change for currents along the $y$ direction in Fig. 2. Conversely when BIA and SIA terms subtract (see Eq. 3) the spin-galvanic effect is suppressed and the CPGE dominates. We would like to emphasize at this point that at the frequency where CPGE is equal to zero for both directions, the current obtained is caused by the spin-galvanic effect only.

The occurrence of a spin-galvanic current is due to the spin dependence of the electron scattering matrix elements $M_{k'k}$. The $2 \times 2$ matrix $\hat{M}_{k'k}$ can be written as a linear combination of the unit matrix $\hat{I}$ and Pauli matrices as follows

$$\hat{M}_{k'k} = A_{k'k}\hat{I} + \sigma \cdot B_{k'k}.$$  

(5)

where $A_{k'k} = A_{kk'}$, $B_{k'k} = B_{kk'}$ due to hermiticity of the interaction and $A_{k',-k} = A_{kk'}$, $B_{k',-k} = -B_{kk'}$ due to the symmetry under time inversion. The spin-dependent part of the scattering amplitude in (001)-grown QW structures is given by

$$\sigma \cdot B_{k'k} = v(k - k')[\sigma_x(k'_y + k_y) - \sigma_y(k'_x + k_x)].$$  

(6)
where \( v(\mathbf{k} - \mathbf{k}') \) is a function defined in [10]. We note that Eq. (6) determines the spin relaxation time, \( \tau_{s}' \), due to the Elliot-Yafet mechanism. Then, for instance, for the spin component \( S_x \) assuming a Boltzmann distribution, the spin-galvanic current in \( y \) direction has the form

\[
j_{SGE,y} = 4\pi e \frac{S_x}{m^* \sum_{kk'} (\tilde{k}_y' - \tilde{k}_y) (\tilde{k}_x' + \tilde{k}_x)^2 |v(\tilde{k} - \tilde{k}' - 2k_0)|^2 \tau_p}
\]

\[
\times f\left(\frac{\hbar^2 \tilde{k}^2}{2m^*}\right) \delta\left(\frac{\hbar^2 \tilde{k}'^2}{2m^*} - \frac{\hbar^2 \tilde{k}^2}{2m^*}\right)
\]

where \( e \) is the electron charge, \( \tau_p \) is the momentum scattering time, \( f \) is the distribution function, \( \delta \) is the delta function, \( m^* \) is the electron effective mass, \( \tilde{k} = k + k_0 \), \( \tilde{k}' = k' - k_0 \), and \( k_0 = (m^* \beta_{xy}/\hbar^2, 0, 0) \). By using Eq. (7) we can estimate the spin-galvanic current as

\[
j_{SGE,y} = Q_{yx} S_x \propto e n_e \frac{\beta_{xy}^{(1)}}{\hbar} \frac{\tau_p}{\tau_{s}'} S_x .
\]

Since scattering is the origin of the spin-galvanic effect, the spin-galvanic current, \( j_{SGE} \), is determined by the Elliot-Yafet spin relaxation time. The relaxation time \( \tau_{s}' \) is proportional to the momentum relaxation time \( \tau_p \). Therefore the ratio \( \tau_p/\tau_{s}' \) in Eq. (8) does not depend on the momentum relaxation time. The in-plane average spin \( S_x \) in Eq. (8) decays with the total spin relaxation time \( \tau_s \) (which may have a contribution from any spin relaxing process). Thus the time decay of the spin-galvanic current following the pulsed photoexcitation is determined by \( \tau_s \). The current in \( x \) direction may be obtained by exchanging \( x \) and \( y \) in Eq. (8).

For the present case, where spin relaxation is obtained as a result of inter-subband absorption of circularly polarized radiation, the current is given by

\[
j_{SGE,x} \sim e \frac{\beta_{yx}^{(1)}}{\hbar} \frac{\tau_p \tau_s}{\hbar \omega} \eta_{21} I \tilde{e}_y ,
\]

\[
j_{SGE,y} \sim e \frac{\beta_{xy}^{(1)}}{\hbar} \frac{\tau_p \tau_s}{\hbar \omega} \eta_{21} I \tilde{e}_x .
\]

\( \eta_{21} \) is the absorbance at the transitions between \( e1 \) and \( e2 \) subbands, \( I \) is the radiation intensity. The parameter \( \xi \) varying between 0 and 1 is the ratio of photoexcited electrons relaxing to the \( e1 \) subband with and without spin-flip. It determines the degree of spin polarization in the lowest subband (see Fig. 3b) and depends on the details of the relaxation mechanism. Optical orientation requires \( \xi \neq 0 \). Eqs. (9) show
that the spin-galvanic current is proportional to the absorbance and is determined by the spin splitting in the first subband, $\beta_{yx}^{(1)}$ or $\beta_{xy}^{(1)}$.

In conclusion we observed the spin-galvanic effect under all-optical excitation and without applying external magnetic fields by making use of the interplay of the Rashba and Dresselhaus splitting of the conduction band. Our results demonstrate in a direct way the non-equivalence of the [110] and [1\bar{1}0] directions in zinc-blende structure QWs. The results also clearly show the difference between the microscopic pictures for spin-galvanic and CPGE, effects which have the same phenomenological description.

We thank L.E. Golub for many helpful discussions. Financial support from the DFG, the RFBR, INTAS, the EPSRC (UK) and FOM (NL) is gratefully acknowledged. The authors are grateful to the FELIX facility team and especially for the skillful assistance of Dr. A.F.G. van der Meer.

[1] Semiconductor Spintronics and Quantum Computation, eds. D.D. Awschalom, D. Loss, and N. Samarth, in the series Nanoscience and technology, eds. K. von Klitzing, H. Sakaki, and R. Wiesendanger (Springer, Berlin, 2002).
[2] Y.A. Bychkov, and E.I. Rashba, Pis’ma ZhETF 39, 66 (1984) [JETP Lett. 39, 78 (1984)].
[3] S.D. Ganichev, E.L. Ivchenko, V.V. Bel’kov, S.A. Tarasenko, M. Sollinger, D. Weiss, W. Wegscheider, and W. Prettl, Nature (London) 417, 153 (2002).
[4] G.M.H. Knippels, X. Yan, A.M. MacLeod, W.A. Gillespie, M. Yasumoto, D. Oepts, and A.F.G. van der Meer, Phys. Rev. Lett. 83, 1578 (1999).
[5] S.D. Ganichev, E.L. Ivchenko, S.N. Danilov, J. Eroms, W. Wegscheider, D. Weiss, and W. Prettl, Phys. Rev. Lett. 86, 4358 (2001).
[6] S.D. Ganichev, V.V. Bel’kov, Petra Schneider, E.L. Ivchenko, S.A. Tarasenko, D. Schuh, W. Wegscheider, D. Weiss, E.V. Beregulin, and W. Prettl, submitted to Phys. Rev. B., cond-mat 0303054.
[7] S.A. Tarasenko, E.L. Ivchenko, V.V. Bel’kov, S.D. Ganichev, D. Schwalter, Petra Schneider, M. Sollinger, and W. Prettl, V.M. Ustinov, A.E. Zhukov, and L.E. Vorob-
[8] M.I. D’yakonov, and V.Yu. Kachorovskii, Fiz. Tekh. Poluprov. 20, 178 (1986) [Sov. Phys. Semicond. 20, 110 (1986)].

[9] O. Krebs, and P. Voisin, Phys. Rev. Lett. 77, 1829 (1996).

[10] N.S. Averkiev, L.E. Golub, and M. Willander, J. Phys.: Condens. Matter 14, R271 (2002).

[11] R.R. Parson, Can. J. Phys. 49, 1850 (1971).

[12] Optical Orientation, F. Meier, and B.P. Zakharchenya, Eds. (Elsevier Science Publ., Amsterdam, 1984).

[13] E.L. Ivchenko, and S.A. Tarasenko, JETP, to be published.