Anisotropic superconducting properties including the upper critical field \( H_{c2} \), thermal activation energy \( U_0 \), and critical current density \( J_c \) are systematically studied in a large \( \text{Ca}_1-x\text{La}_x\text{Fe}_2\text{As}_2 \) single crystal \((x \approx 0.18)\). The obtained \( H_{c2} \) bears a moderate anisotropy \( \gamma \) of approximately 2–4.2, located between those of \( \text{Ba}_x\text{K}_2\text{Fe}_2\text{As}_2 \) \((1 < \gamma < 2)\) and \( \text{NdFeAsO}_{1-x}\text{F}_x \) \((5 < \gamma < 9.2)\). Both the magnitude of \( U_0 \) and its field dependence are very similar to those of \( \text{NdFeAsO}_{1-x}\text{F}_x \), also indicating anisotropic superconductivity. Moreover, high and anisotropic \( J_c \)'s exceeding \( 10^5 \text{A/cm}^2 \) have been calculated from the magnetization hysteresis loops, indicating the existence of strong bulk-dominated pinning in the present superconducting material. © 2014 The Japan Society of Applied Physics
the two-band theory.\textsuperscript{15} Nevertheless, the linear $T$ dependencies of $H_{c2}$ above 1 T satisfy well the Werthamer–Helfand–Hohenberg (WHH) model.\textsuperscript{16} According to the WHH model, the upper critical field at zero temperature is roughly estimated to be $H_{c2}^{(0)} = 39.4$ T ($H \parallel c$, $H_{c2}$ slope = 14 kOe/K) and $H_{c2}^{(0)} = 166.2$ T ($H \parallel ab$, $H_{c2}$ slope = 59 kOe/K). The corresponding coherence lengths are $\xi_c (0) = 6.9$ Å and $\xi_{ab} (0) = 28.9$ Å. The $H_{c2}$ anisotropy at zero temperature $\gamma (0)$ is 4.2. It should be pointed out that the estimated $H_{c2} (0)$, $\xi (0)$, and $\gamma (0)$ should be modified for the possible existence of anisotropic paramagnetic effects as in many IBSs.\textsuperscript{7,17}

Therefore, we use the $\gamma$ value ($\sim 2.08$) near $T_c$ for comparison with the other IBSs. On one hand, such a $\gamma$ value locates between those of “122”-type Ba$_{1-\delta}$K$_{\delta}$Fe$_2$As$_2$ ($1 < \gamma < 2$)\textsuperscript{7} and “1111”-type NdFeAsO$_{1-x}$F$_x$ ($5 < \gamma < 9.2$).\textsuperscript{18,19} As mentioned above, $\gamma$ reflects directly the coupling strength between the charge reservoir layers and the conducting FeAs layers. That is, the distance $d$ between the adjacent FeAs layers plays an important role in determining $\gamma$ values. A large distance $d$ may lead to large $\gamma$ values. The $d$ values of Ca$_{1-x}$La$_x$FeAs$_2$ ($c = 10.35$ Å) and NdFeAsO$_{1-x}$F$_x$ ($c = 8.56$ Å)\textsuperscript{10} are the same as the $c$-axis lattice constant $c$, while the formula $d = c/2$ should be used for Ba$_{1-\delta}$K$_{\delta}$Fe$_2$As$_2$ ($c = 13.30$ Å)\textsuperscript{19} owing to the existence of double FeAs planes in a single unit cell. Accordingly, Ba$_{1-\delta}$K$_{\delta}$Fe$_2$As$_2$ exhibits the smallest anisotropy. Compared with Ca$_{1-x}$La$_x$FeAs$_2$, “1111”-type NdFeAsO$_{1-x}$F$_x$ bears a smaller FeAs distance $d$ but exhibits a larger $\gamma$. Possibly, the $\gamma$ values may also be affected by other factors, such as the electrical conductivity of the reservoir layers [Nd(O,F)$_x$(Ba,K)$_{1-x}$, and (Ca,La,As)$_{1-x}$]. Here, $x$ means the valence] stacked between FeAs planes. Because the reservoir layer Nd(O,F)$_x$ in NdFeAsO$_{1-x}$F$_x$ bears poor electrical conductivity, a slightly large anisotropy is reflected. On the other hand, compared with the extremely large $H_{c2}$ anisotropy in the “122”-type analogy superconductors Ca$_{1-x}$RE$_x$Fe$_2$As$_2$ whose weak high-$T_c$ superconductivity is of filamentary nature,\textsuperscript{20,22} $\gamma$ in Ca$_{1-x}$La$_x$FeAs$_2$ is clearly small. In this view, the superconductivity origin in the new “112”-type compounds may be different from that of the filamentary superconductivity in Ca$_{1-x}$RE$_x$Fe$_2$As$_2$.

This speculation is also evidenced by the magnetic results discussed later. In the bottom-left inset of Fig. 2, temperature-dependent $\gamma$ near $T_c$ is shown. Similar to the trend in NdFeAsO$_{1-x}$F$_x$, $\gamma$ increases gradually from 2 to a value above 4 with decreasing temperature. This trend is opposite to that of Ba$_{1-\delta}$K$_{\delta}$Fe$_2$As$_2$, which may be caused by multiband effects.\textsuperscript{7}

Generally, for the resistive transition in the presence of a magnetic field, the effect of thermally activated flow of vortices should be taken into account, which means that the temperature dependence of resistivity in this region is expected to be expressed by the Arrhenius plot,\textsuperscript{19} $\rho (T,H) = \rho_0 \exp [–U(T,H)/k_B T]$. Here, $U(T,H)$ is the activation energy of the flux flow (also known as the pinning potential/energy), $\rho_0$ is a constant, and $k_B$ is Boltzmann’s constant. As can be seen in the upper-right and bottom-left insets of Fig. 3, the Arrhenius dependence holds about 2–3 decades. This linear behavior guarantees the suitable application of the formula $\rho (T,H) = \rho_0 \exp [–U_0(H)/k_B T]$.\textsuperscript{19,23} The activation energy $U_0(H)$ (main panel of Fig. 3) thus can be extracted from the slope of the curve $\log \rho (T,H)$ versus $1/T$. As can be seen, $U_0$ for both field orientations is approximately $(3–4) \times 10^3$ K at low fields. The field dependencies of $U_0$ for $H \parallel ab$ and $H \parallel c$ are markedly different. For $H \parallel c$, $U_0$ shows a very weak dependence in the low-magnetic-field region ($H < 1$ T). A stronger power law decrease $U_0 \propto B^{-0.56}$ takes place for $H$ above 1 T. A similar feature of $U_0$ with double field dependences was also found in cuprates and NdFeAsO$_{1-x}$F$_x$,\textsuperscript{19} which was explained by a transition from a single-vortex-dominated pinning to a small bundle pinning. In the state of the single-vortex-dominated pinning, the overlap of vortices is negligible and the field dependence is weak. When $H$ is increased to a certain amount (noted as $H^*$) that results in significant overlaps of vortices, the pinning energy $U_0$ begins to be markedly suppressed. For $H \parallel ab$, a single power law $U_0 \propto B^{-0.29}$ holds through the field regions, which possibly results from a nearly negligible $H^*$. It is peculiarly interesting that, similar to NdFeAsO$_{1-x}$F$_x$, a slightly large $U_0$ is found for $H \parallel c$ in the low-field region in contrast to the result of Ba$_{1-\delta}$K$_{\delta}$Fe$_2$As$_2$.\textsuperscript{24} We believe that the larger $U_0$ for $H \parallel c$
in Ca$_{1-x}$La$_x$FeAs$_2$ and NdFeAsO$_{1-y}$F$_y$ is caused by the anisotropic $H^*$, which is mainly determined by the anisotropic superconducting penetration depth $\lambda$ ($\lambda_0 < \lambda_c$). The large $\lambda_c$ leads to a nearly negligible $H^*$ for $H \parallel ab$. In the field region $H^*(H \parallel ab) < H < H^*(H \parallel c)$, the single-vortex pinning is dominated for $H \parallel c$, while the vortex pinning for $H \parallel ab$ has already stepped into the small bundle state. The weaker field dependence of $U_0$ in the single-vortex pinning state causes the larger value of $U_0$ ($H \parallel c$). However, in the high-field region, since $H_{c2}(H \parallel ab)$ is larger than $H_{c2}(H \parallel c)$, $U_0(H \parallel c)$ tends to decrease faster with the field than $U_0(H \parallel ab)$, resulting in a crossover of $U_0$. Overall, the anisotropic behavior of $U_0$ reflects a relatively two-dimensional superconductivity compared with that of Ba$_{1-x}$K$_x$Fe$_2$As$_2$.

Magnetic hysteresis loops (MHLs) at 5 K with fields applied both parallel to the $c$-axis ($H \parallel c$) and the $ab$-plane ($H \parallel ab$) were also collected [Fig. 4(a)]. The large and symmetric MHLs imply that the bulk pinning dominates in the crystal. Such large MHLs are observed for the first time in the present superconducting material. For a rectangular prism-shaped crystal with dimensions of $2c < 2a < 2b$, there are three types of critical current $J_{c}^{x,y}$, where $x$ and $y$ refer to the directions of the current and magnetic field, respectively. When the magnetic field is applied along the $c$ side, the supercurrent density $J_{c}^{ab,c}$ is generated by the gradients of the Abrikosov vortices perpendicular to the $ab$-plane, which can be determined easily using the Bean model. Figure 4(b) shows the calculated critical current density $J_{c}^{ab,c}$ under the Bean construction,

$$J = \frac{20\Delta M}{a(1-a/3b)},$$

where $\Delta M$ (unit: emu/cm$^3$) is $M_{down} - M_{up}$; $M_{down}$ and $M_{up}$ are magnetizations when sweeping fields down and up, respectively. Here, $a$ and $b$ (unit: cm) are the redefined sample width and length ($a < b$) under a certain configuration of field orientation. As can be seen, the self-field $J_{c}^{ab,c}$ at 5 K reaches a high value of $3.6 \times 10^5$ A/cm$^2$. Such a value is comparable to that of a well-annealed “11” crystal Fe$_{1+y}$(Te,Se)$_8$ whose magnetic critical current density is almost isotropic.

On the other hand, when the magnetic field is applied along the $a$ (or $b$) side, the situation becomes more complicated for the existence of two current densities with different directions, of which one is along the $ab$-plane ($J_{c}^{ab,ab}$) and another is across the planes ($J_{c}^{ab,ab}$). Assuming that $J_{c}^{ab,ab}$ is equal to $J_{c}^{ab,ab}$, the weighted average $J_{c}^{mean}$ for $H \parallel ab$ can be directly obtained using Eq. (1). The $J_{c}^{mean}$ values clearly exceed $10^5$ A/cm$^2$. Even though $J_{c}^{mean}$ is frequently used in the $H \parallel ab$ case in many superconductors, the above assumption is somewhat unreasonable. By the method described in Ref. 26 and assuming $J_{c}^{ab,ab} = J_{c}^{ab,ab}$, relatively reliable $J_{c}^{ab,ab}$ values are obtained, as shown in Fig. 4(b). $J_{c}^{ab,ab}$ reflects much smaller values than $J_{c}^{mean}$, which is more reasonable for the reason that the coherence lengths $\xi_c$ (< 6.9 Å) are smaller than the FeAs distance $d$ (≈10.35 Å). Therefore, similar to the $H_{c2}$ anisotropy $\gamma$, a larger $J_{c}^{ab,ab}$ is found for Ca$_{1-x}$La$_x$FeAs$_2$ than for the nearly isotropic “122” compounds. It should also be noted here that for the magnetic field applied both parallel to the $c$-axis and the $ab$-plane, the values of $J_{c}^{ab,ab}$ (5 K) slightly change against magnetic fields up to 6 T, revealing a magnetic-field-robust feature. The large and magnetic-field-robust $J_{c}$ could not originate from interface or filamentary superconductivity. Actually, in the case of the filamentary superconductivity in Ca$_{1-x}$RE$_x$Fe$_2$As$_2$, almost no superconducting MHLs can be measured.

Therefore, in terms of $T_c$ and superconductivity anisotropy, the “112”-type Ca$_{1-x}$La$_x$FeAs$_2$ ($T_c \sim 42.6$ K) exhibits relatively moderate values, bridging the difference between “122” ($T_c \sim 38$ K) and “1111” ($T_c \sim 55$ K) iron-based compounds. Possibly, similar to cuprates, a high anisotropy may act as an important ingredient for the achievement of high transition temperatures.

In summary, large Ca$_{0.77}$La$_{0.18}$Fe$_{0.9}$As$_2$ single crystals with the best superconducting performance are synthesized. The roughly estimated upper critical fields $H_{c2}(0)$’s using the WHH model in the present superconducting material reach values as high as 39.4 and 166.2 T for the out-plane ($H \parallel c$) and in-plane ($H \parallel ab$) directions, respectively. The $H_{c2}$ anisotropy $\gamma$ near $T_c$ is 2.08. The anisotropic pinning potential $U_0$ shows comparable behaviors to that of “1111” NdFeAsO$_{1-y}$F$_y$. A high critical current density $J_{c}^{xx}$ of over $10^5$ A/cm$^2$ is determined from magnetization hysteresis loops, indicating strong bulk-dominated pinning. The moderate anisotropy in the new “112”-type superconductor seems to bridge the difference between “122”- and “1111”-type IBSs.
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1) N. Katayama, K. Kudo, S. Onari, T. Mizukami, K. Sugawara, Y. Sugiyama, Y. Kitahama, K. Iba, K. Fujimura, N. Nishimoto, M. Nohara, and H. Sawa, J. Phys. Soc. Jpn. 82, 123702 (2013).
2) H. Yakita, H. Ogino, T. Okada, A. Yamamoto, K. Kishio, T. Tohei, Y. Ikushara, Y. Gotoh, H. Fujihisa, K. Kataoka, H. Eisaki, and J. Shimoyama, J. Am. Chem. Soc. 136, 846 (2014).
3) Y. Kamihara, T. Watanabe, M. Hirano, and H. Hosono, J. Am. Chem. Soc. 130, 3296 (2008).
4) M. Rotter, M. Tegel, and D. Johrendt, Phys. Rev. Lett. 101, 107006 (2008).
5) N. Ni, S. L. Bud’ko, A. Kreyssig, S. Nandi, G. E. Rustan, A. I. Goldman, S. Gupta, J. D. Corbett, A. Kracher, and P. C. Canfield, Phys. Rev. B 78, 014507 (2008).
6) H. Ding, P. Richard, K. Nakayama, K. Sugawara, T. Arakane, Y. Sekiba, A. Takayama, S. Souma, T. Sato, T. Takahashi, Z. Wang, X. Dai, Z. Fang, G. F. Chen, J. L. Luo, and N. L. Wang, Europhys. Lett. 83, 47001 (2008).
7) H. Q. Yuan, J. Singleton, F. F. Balakirev, S. A. Baily, G. F. Chen, J. L. Luo, and N. L. Wang, Phys. Rev. B 81, 020509 (2010).
8) Y. Jia, P. Cheng, L. Fang, H. Luo, H. Yang, C. Ren, L. Shan, C. Gu, and H.-H. Wen, Appl. Phys. Lett. 93, 032503 (2008).
9) J. Jaroszynski, F. Hunte, I. Raičević, A. Gurevich, D. C. Larbalestier, F. F. Balakirev, L. Fang, P. Cheng, Y. Jia, and H. H. Wen, Phys. Rev. B 78, 174523 (2008).
10) B. Lv, L. Deng, M. Gooch, F. Wei, Y. Sun, J. K. Meen, Y.-Y. Xue, B. Lorenz, and C.-W. Chu, Proc. Natl. Acad. Sci. U.S.A. 108, 15705 (2011).
11) W. Zhou, F. F. Yuan, J. C. Zhuang, Y. Sun, Y. Ding, L. J. Cui, J. Bai, and Z. X. Shi, Supercond. Sci. Technol. 26, 095003 (2013).
12) T. Tamegai, Q. P. Ding, T. Ishibashi, and Y. Nakajima, Physica C 484, 31 (2013).